Neutrino Assisted Gauge Mediation

Hyung Do Kim, Doh Young Mo, Min-Seok Seo

Department of Physics and Astronomy and Center for Theoretical Physics,

Seoul National University, Seoul 151-747, Korea

Abstract

Recent observation shows that the Higgs mass is at around 125 GeV while the prediction of the minimal supersymmetric standard model is below 120 GeV for stop mass lighter than 2 TeV unless the top squark has a maximal mixing. We consider the right-handed neutrino supermultiplets as messengers in addition to the usual gauge mediation to obtain sizeable tri-linear soft parameters A_t needed for the maximal stop mixing. Neutrino messengers can explain the observed Higgs mass for stop mass around 1 TeV. Neutrino assistance can also generate charged lepton flavor violation including $\mu \to e\gamma$ as a possible signature of the neutrino messengers. We consider S_4 discrete flavor model and show the relation of the charged lepton flavor violation, θ_{13} of neutrino oscillation and muon $g-2$.

I. INTRODUCTION

The observation of the Standard Model Higgs-like new boson with mass at around 125 GeV [1, 2] changes the current understanding of new physics at the weak scale. The minimal supersymmetric standard model (MSSM) can explain 125 GeV with a relatively light stop of 1 to 2 TeV in the context of maximal stop mixing. From the model building point of view, it is quite difficult to realise the maximal stop mixing scenario starting from ultraviolet (UV) theory. In minimal gauge mediation (MGM) [3–6], soft tri-linear A term is not generated at the messenger scale and the radiatively generated A term at the weak scale is not large enough to realise the maximal stop mixing. As a result, colored superpartners should be as heavy as 5 to 10 TeV to explain 125 GeV mass of the Higgs boson [7, 8]. Therefore, the explanation of 125 GeV Higgs boson mass needs an extra help in minimal gauge mediation. Next to the minimal supersymmetric standard model (NMSSM) can use the extra contribution from the Yukawa F-term of the singlet. For instance, look at [9]. Extra vector-like fermions are added in minimal gauge mediation [10–12]. Direct coupling of visible sector fields with messengers can help. Higgs-messenger mixing [13, 14] and matter-messenger mixing [15–19] can generate Yukawa mediated contribution including A term at the messenger scale. However, at the same time the virtue of gauge mediation is gone and they would spoil nice flavor preserving spectrum and can possibly cause the flavor problem at the weak scale. For the Higgs-messenger mixing, A/m^2 problem[13, 14] which is analogous to $\mu/B\mu$ problem can arise and the electroweak symmetry breaking is difficult to achieve if the mixing coupling is large. General gauge mediation [20] can avoid this problem by using the mechanism of radiatively generated maximal stop mixing [21, 22].

In this paper we consider the right-handed neutrino supermultiplets as the messengers of supersymmetry breaking in addition to the messengers charged under the Standard Model (SM) gauge group, e.g. $\mathbf{5}$ and $\bar{\mathbf{5}}$ of $SU(5)$. The setup is motivated from [23] which provides a solution to μ problem in gauge mediation (more precisely $\mu/B\mu$ problem) [24, 25]. For the solution in [23] to work, the messenger scale should be higher than the Peccei-Quinn breaking scale, $10^9 \sim 10^{11}$ GeV. For the gauge mediation to be the dominant contribution compared to the Planck suppressed higher dimensional contribution, the messenger scale should be lower than 10^{15} GeV. Therefore, the See-Saw scale with order one neutrino Dirac Yukawa couplings, $10^{13} \sim 10^{14}$ GeV, is well motivated as the messenger scale if we accept

[23] as a solution of the μ problem in gauge mediation. If the messenger scale is at around the See-Saw scale, the natural question is why the right-handed neutrino supermultiplets do not serve as messengers of supersymmetry breaking. Apparently there is no harm to couple the right-handed neutrino superfields directly to the messengers. Majorana mass of the right-handed neutrino and the messenger mass of ordinary $\mathbf{5}$ and $\bar{\mathbf{5}}$ might have the same origin in this case. In summary, the minimal set of messengers are $\mathbf{5}$, $\bar{\mathbf{5}}$ and $\mathbf{1}$. This is different from previous studies relating gauge mediation and See-Saw mechanism [26–28]. They employ particles relevant to See-Saw mechanism as messengers, and these particles are also charged under the SM gauge group, which can be seen in the Type-II or Type-III See-Saw. Therefore, gauge mediation and neutrino Dirac Yukawa mediation have a common messenger. In our case, in contrast, neutrino Dirac Yukawa messenger is the right-handed neutrino, the SM singlet.

If the right-handed neutrinos couple to the supersymmetry breaking field, neutrino Dirac Yukawa coupling generates the A term and soft scalar mass of lepton doublet and up-type Higgs at the See-Saw scale after integrating out the neutrino messengers. 125 GeV Higgs mass can be explained with stop lighter than 2 TeV in this setup. At the same time the stop mass gets an extra Yukawa mediation and maximal stop mixing can be easily realised.

As there is a neutrino Dirac Yukawa contribution to the soft parameters in addition to the ordinary gauge mediation, interesting new physics signature is expected. The mechanism of the charged lepton flavor violation is different from that in mSUGRA [29] or SUSY GUT [30] in which the origin is the running of soft parameters above the See-Saw scale. Though the origin is different, the spectrum looks similar. The crucial difference is that here the flavor violation appearing in lepton doublet soft scalar mass is $16\pi^2$ bigger than the one in mSUGRA. Therefore, the naive expectation is that order one neutrino Dirac Yukawa coupling would be incompatible with the current bounds of various charged lepton flavor violation constraints including $\mu \to e\gamma$.

The computation of the charged lepton flavor violation needs a complete flavor model. Current observation of the charged lepton mass and lepton mixing matrix (PMNS) can be explained in a consistent way with the neutrino Dirac Yukawa matrix which is proportional to the identity matrix. This is not an ad hoc assumption but can be explained in the context of non-Abelian discrete flavor symmetry, e.g., tribimaximal PMNS[31] from S_4. Therefore, order one neutrino Dirac Yukawa coupling can generate order one A term at

the messenger scale and at the same time can be consistent with the charged lepton flavor violation constraints as long as it is proportional to the identity matrix.

S_4 flavor symmetry is the most natural and/or simple if $\theta_{13} = 0$ as the tribimaximal mixing can be nicely realised. However, small but sizeable θ_{13} ($\sin\theta_{13} \sim 0.15$) can be accommodated with the extra complication[32–35]. If the origin of θ_{13} is the modification of Majorana mass of the right-handed neutrino, there would be no off-diagonal element in the lepton doublet soft scalar masses as the neutrino Dirac Yukawa would be still proportional to the identity. In this case the model is free from the cLFV constraints. Nevertheless, the sparticle spectrum needed to explain the observed Higgs mass is heavy enough such that it is hard to explain the muon anomalous magnetic moment at the same time. If θ_{13} is due to the deviation of the neutrino Dirac Yukawa matrix from the identity, sizeable charged lepton flavor violation is expected. We compute the charged lepton flavor violating processes in both cases and show that interesting parameter space exists if θ_{13} is a combination of two contributions from neutrino Dirac Yukawa and Majorana mass matrix.

The contents of the paper is following. In section 2, we explain the setup for neutrino assisted gauge mediation in which the right-handed neutrino is added as messengers in addition to the ordinary SM charged messengers. Also we discuss the implication for the Higgs mass. In section 3, we explain our S_4 flavor model as a representative example to discuss possible phenomenological implication. In section 4, we discuss charged lepton flavor violation in connection with muon anomalous magnetic moment, the neutrino mixing angle θ_{13} and the Higgs mass. Then we conclude.

II. NEUTRINO ASSISTED GAUGE MEDIATION AND THE HIGGS MASS

A. Soft terms generated from right-handed neutrino messengers

The extremely small masses of neutrinos can be explained through the See-Saw mechanism[36–40], in which lepton number is violated at around the Grand Unified Theory (GUT) scale. In this paper, we consider the simplest model, type-I See-Saw. For this, we extend the MSSM superpotential by including right-handed Majorana neutrinos,

$$\begin{aligned} W = \epsilon_{ab}\Big[&(Y_U)_{ij}\bar{U}_i Q^a_j H^b_u - (Y_D)_{ij}\bar{D}_i Q^a_j H^b_d - (Y_E)_{ij}\bar{E}_i L^a_j H^b_d + (Y_\nu)_{ij} N_i L^a_j H^b_u \\ &+ \mu H^a_u H^b_d\Big] + \frac{1}{2} M^{ij}_N N_i N_j, \end{aligned} \tag{1}$$

where ϵ_{ab} is a totally antisymmetric tensor with $\epsilon_{12} = 1$. The superfields in the superpotential represent right-handed neutrino-sneutrino pairs, in addition to the SM particles and their superpartners. They have the following SM gauge group $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ quantum numbers:

$$\begin{aligned} & Q : (3, 2, \frac{1}{6}), \quad \bar{U} : (\bar{3}, 1, -\frac{2}{3}), \quad \bar{D} : (\bar{3}, 1, \frac{1}{3}) \\ & L : (1, 2, -\frac{1}{2}), \quad \bar{E} : (1, 1, 1), \quad N : (1, 0, 0) \\ & H_u : (1, 2, \frac{1}{2}), \quad H_d : (1, 2, -\frac{1}{2}). \end{aligned} \tag{2}$$

Relative minus signs of Yukawa terms are given to make the sign of terms responsible for the fermion Dirac masses to be the same.

The relevant soft supersymmetry (SUSY) breaking terms are given by

$$\begin{aligned} \mathcal{L}_{\rm soft} = & -(m_N^2)^i_j \tilde{N}^\dagger_i \tilde{N}_j - (m_L^2)^j_i \tilde{L}^{\dagger i} \tilde{L}_j - m_{H_u}^2 H_u^\dagger H_u \\ & - \Big[\frac{1}{2} (B_N M)^{ij} \tilde{N}_i \tilde{N}_j + (\tilde{A_U})_{ij} \tilde{U}^i \tilde{Q}^j H_u - (\tilde{A_D})_{ij} \tilde{D}^i \tilde{Q}^j H_d - (\tilde{A_E})_{ij} \tilde{E}^i \tilde{L}^j H_d + B\mu H_u H_d + h.c. \Big]. \end{aligned} \tag{3}$$

We consider two origins of soft terms. The first one is gauge mediation. In the gauge mediation, sfermions obtain soft masses given by [6]

$$m_{\tilde{f}}^2 = 4 \sum_a \left(\frac{g_a^2}{16\pi^2} \right)^2 C_a \sum_i \left(\frac{F}{M_i} \right)^2 T_a(\mathcal{R}_i) f(x_i) \tag{4}$$

at the messenger scales M_i, where C_a is the quadratic Casmir $\sum_\alpha T^\alpha T^\alpha$ of the sfermion representation $\mathcal{R}_i$ under the corresponding gauge group labeled by a, which is given by $(N^2 - 1)/(2N)$ for SU(N) and Y^2 for $\mathrm{U}(1)_Y$, T_a is defined by $\mathrm{Tr} T^\alpha T^\beta = T_a(\mathcal{R}_i) \delta^{\alpha\beta}$, and $f(x_i)$ is the loop function of $x_i = F/M_i^2$ which is close to one for small x_i. On the other hand, the 2-loop tri-linear A term is very small and can be neglected at the messenger scale.

For the second origin of soft terms, we introduce a SUSY breaking spurion X which couples to the right-handed neutrinos. Majorana mass of the right-handed neutrino comes from the scalar vacuum expectation value (VEV) of the SUSY breaking spurion X,

$$W \supset \lambda X N N. \tag{5}$$

Then N acts as the messengers of supersymmetry breaking, and the neutrino Dirac Yukawa coupling,

$$W \supset Y_\nu N L H_u, \tag{6}$$

is interpreted as the direct mixing term among the messengers, Higgs and matter (leptons).

The SUSY breaking effects at the See-Saw scale $M_N = \lambda\langle X\rangle$ is studied in [41]. When right-handed neutrinos couple to the SUSY breaking sector, Majorana mass matrix is analytically continued to be $M_N \to (1+\theta^2 B_N)M_N$, as in the case of gauge mediation[42–44]. Here, we assume that the flavor structure of the right-handed neutrinos is fully determined by M_N, so $B_N = F_X/X$ is a constant.

Then, SUSY breaking is transferred to the visible sector through the neutrino Dirac Yukawa interaction. Wave function renormalization from the interaction with right-handed neutrinos is given by

$$\delta Z_L = \frac{Y_\nu^{R\dagger}}{16\pi^2}\Big(1-\ln\frac{M^{R\dagger}M^R}{\Lambda^2}\Big)Y_\nu^R, \qquad \delta Z_{H_u} = \mathrm{Tr}\delta Z_L \tag{7}$$

where

$$\lambda_N^R = [Z_N^{-1/2}]^T\lambda_N Z_L^{-1/2}Z_{H_u}^{-1/2}, \qquad M^R = [Z_N^{-1/2}]^T M_N Z_N^{-1/2}, \tag{8}$$

then analytically continued Majorana masses give the soft masses. From field redefinitions

$$\begin{aligned} L &\to \big(1-\frac{\delta Z_L|_0}{2}\big)(1-\theta^2\delta Z_L|_{\theta^2})L \\ H_u &\to \big(1-\frac{\delta Z_{H_u}|_0}{2}\big)(1-\theta^2\delta Z_{H_u}|_{\theta^2})H_u, \end{aligned} \tag{9}$$

supersymmetric kinetic terms can be written in the simple form,

$$\Phi^\dagger(1+\delta Z_\Phi)\Phi \to \Phi^\dagger(1+\theta^2\bar\theta^2\delta Z_\Phi|_{\theta^2\bar\theta^2})\Phi \tag{10}$$

then we can read off the one-loop corrections to the soft masses

$$\delta m_L^2 = -\delta Z_L|_{\theta^2\bar\theta^2} \quad \text{and} \quad \delta m_{H_u}^2 = -\delta Z_{H_u}|_{\theta^2\bar\theta^2}. \tag{11}$$

In the expression, B_N is just a constant, not a matrix. So $\ln(M_N^\dagger M_N)$ in the wave function renormalization is separated into holomorphic and anti-holomorphic parts, respectively. Since $\theta^2\bar\theta^2$ term is not generated, we do not have one-loop soft masses.

Hence, as in minimal gauge mediation, soft masses are generated at two loop level. In [13], it was shown that soft scalar masses of the fields which directly couple to messengers and those which do not are different. In our model, the slepton $\tilde L$ and the up-type Higgs H_u couple to messengers N directly to give soft terms,

$$
\begin{aligned}
\delta m_L^2 &= \frac{B_N^2}{(4\pi)^4}\Big[\Big(\mathrm{Tr}[Y_\nu Y_\nu^\dagger] + 3\mathrm{Tr}[Y_U Y_U^\dagger] - 3g_2^2 - \frac{1}{5}g_1^2\Big)Y_\nu^\dagger Y_\nu + 3Y_\nu^\dagger Y_\nu Y_\nu^\dagger Y_\nu\Big] \\
\delta m_{H_u}^2 &= \frac{B_N^2}{(4\pi)^4}\Big[4\mathrm{Tr}[Y_\nu Y_\nu^\dagger Y_\nu^\dagger Y_\nu] - \Big(3g_2^2 + \frac{1}{5}g_1^2\Big)\mathrm{Tr}[Y_\nu Y_\nu^\dagger]\Big].
\end{aligned} \tag{12}
$$

On the other hand, $\tilde{Q}$ and $\tilde{U}$ obtain two-loop soft scalar masses through the wave function renormalization of H_u and the corrections are given by

$$
\begin{aligned}
\delta m_Q^2 &= -\frac{B_N^2}{(4\pi)^4}\mathrm{Tr}[Y_\nu Y_\nu^\dagger]Y_U^\dagger Y_U \\
\delta m_U^2 &= -\frac{B_N^2}{(4\pi)^4}\mathrm{Tr}[Y_\nu Y_\nu^\dagger]Y_U Y_U^\dagger
\end{aligned} \tag{13}
$$

while the soft masses of $\tilde{E}$ and H_d come out of the wave function renormalization of L and the corrections are given by

$$
\begin{aligned}
\delta m_E^2 &= -\frac{B_N^2}{(4\pi)^4}Y_E Y_\nu^\dagger Y_\nu Y_E^\dagger \\
\delta m_{H_d}^2 &= -\frac{B_N^2}{(4\pi)^4}\mathrm{Tr}[Y_E Y_\nu^\dagger Y_\nu Y_E^\dagger].
\end{aligned} \tag{14}
$$

By replacing $Y_E \to Y_E(1+\delta A_E)$, $Y_U \to Y_U(1+\delta A_U)$, and $Y_D \to Y_D(1+\delta A_D)$, we have following soft terms at one loop level,

$$
\begin{aligned}
&\delta A_E = -\delta Z_L|_{\theta^2}, \quad \delta A_U = -\mathbb{I}\delta Z_{H_u}|_{\theta^2}, \\
&\delta A_D = 0, \quad \delta B = -\delta Z_{H_u}|_{\theta^2}.
\end{aligned} \tag{15}
$$

Unlike gauge mediation, right-handed neutrino mediation generates one-loop $A-$terms,

$$
\begin{aligned}
A_E &= \frac{B_N}{16\pi^2}Y_\nu^\dagger Y_\nu \\
A_U &= -\mathrm{Tr}A_E \times \mathbb{I}_{3\times 3} \\
B &= \mathrm{Tr}A_E.
\end{aligned} \tag{16}
$$

While gauge mediation contributions are flavor universal, See-Saw Yukawa mediation is flavor dependent and one of the virtue of the gauge mediaion would disappear. In the absence of See-Saw Yukawa mediation, cLFV can appear when the messenger scale is higher than the right-handed neutrino Majorana mass scale. See-Saw Yukawa contributes to the slepton soft mass through the renormalization group equation (RGE) ,

$$
\mu\frac{d}{d\mu}m_L^2 = \mu\frac{d}{d\mu}m_L^2\Big|_{\mathrm{MGM}} + \frac{1}{16\pi^2}\Big[(m_L^2 Y_\nu^\dagger Y_\nu + Y_\nu^\dagger Y_\nu m_L^2) + 2(Y_\nu^\dagger m_N^2 Y_\nu + m_{H_u}^2 Y_\nu^\dagger Y_\nu + \tilde{A}_\nu^\dagger \tilde{A}_\nu)\Big] \tag{17}
$$

which should be restricted by cLFV constraints [45]. Here $\tilde{A}_\nu = A_\nu Y_\nu$ is used. Since m_L^2 is two-loop generated, cLFV effects are further loop suppressed (at three loop level). Unlike mSUGRA, this effect is known to be small in gauge mediation as the messenger scale is at most comparable to the See-Saw scale and the running can be made in a very short interval. This is not the cLFV that we are interested in.

In neutrino assisted gauge mediation, neutrino Dirac Yukawa couplings can introduce two-loop generated cLFV effects on m_L^2 as a result of gauge-Yukawa or Yukawa mediation,

$$\delta m_L^2 = \frac{B_N^2}{(4\pi)^4}\Big[\Big(\mathrm{Tr}[Y_\nu Y_\nu^\dagger] + 3\mathrm{Tr}[Y_U Y_U^\dagger] - 3g_2^2 - \frac{1}{5}g_1^2\Big)Y_\nu^\dagger Y_\nu + 3Y_\nu^\dagger Y_\nu Y_\nu^\dagger Y_\nu\Big] \tag{18}$$

in the charged lepton mass basis. If the two loop generated slepton mass squared has a nonzero off-diagonal element, it would generate cLFV. Parametrically, this effect is much larger than the expected cLFV in mSUGRA or similar scenarios in which the effect comes from the running above the See-Saw scale. We simply assume that both messengers $\mathbf{5}, \bar{\mathbf{5}}$ and $\mathbf{1}$ have the same masses at the See-Saw scale. In principle these two masses can be different and cLFV can arise if the singlet messenger is lighter than $\mathbf{5}, \bar{\mathbf{5}}$. However, this effect is loop suppressed compared to the Yukawa mediation we would not consider it in this paper.

Further discussion on cLFV is possible only when there is an explicit flavor model providing the neutrino Dirac Yukawa and charged lepton Yukawa matrices. As a simple and illustrative example of the explicit model, we consider S_4 flavor symmetry in Sec. III. It will be shown that various types of See-Saw Yukawa Y_ν would predict different sizes of effects on cLFV. Before moving onto the flavor discussion, let us consider the implication on the Higgs mass first.

B. Higgs mass and superparticle spectrum

Minimal gauge mediation does not generate A_t at one loop and the weak scale A_t is radiatively generated by the gluino loop. However, the same gluino contribution appears in stop soft scalar mass and the relative ratio of $|A_t|$ and $m_{\tilde{t}}$ can not be large. On the other hand, the physical light CP even Higgs mass in the MSSM is affected by $\hat{X}_t \equiv (A_t - \mu/\tan\beta)/m_{\tilde{t}}$ and $\hat{X}_t \sim 2$ (or $\sqrt{6}$ more precisely) gives the maximum finite threshold correction as shown in Fig. 1.

One way to make $|\hat{X}_t| > 1$ at the weak scale is to start from tachyonic stop boundary

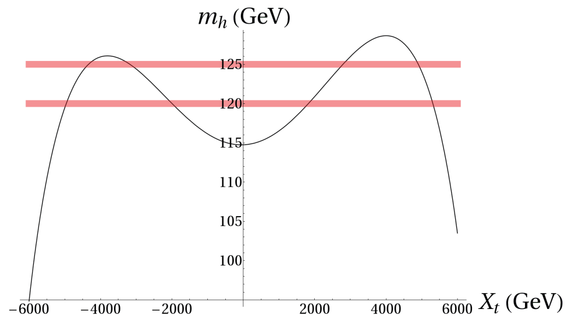

FIG. 1. Higgs mass with respect to X_t for $\tan\beta = 10$, stop mass $M_{\tilde{t}} \sim 2\,\text{TeV}$.

condition [21] which is explicitly realised in gauge messenger model [22]. However, this option is not available in minimal gauge mediation. The other possibility is to couple messengers directly to the visible sector fields such that large negative A term can be generated at the messenger scale. If A term is positive, the gluino contribution from the running cancels the A term at the messenger scale. Matter-messenger mixing [15–19] also has been considered recently. Messenger-matter-matter Yukawa coupling would generate the needed A_t term at the messenger scale. However, the full Yukawa couplings are written as 3×3 matrices and why all other dangerous Yukawa couplings between matters and messengers are absent except 33 component remains to be a puzzle. One way out is to consider Higgs-messenger mass mixing [44] and to generate all the wanted Yukawa couplings between matter and messengers from ordinary Yukawa couplings of matter with Higgs. There would be a direct modification of squark spectrum if squark couples directly to the messenger.

Higgs-messenger mixing through Higgs-messenger-messenger coupling or Higgs-Higgs-messenger coupling has been considered in [13, 14]. In this case, we often encounter A/m^2 problem. To understand this, it is worth to emphasize that, the two-loop soft mass squared of the Higgs field H_u which has direct coupling to messenger Φ has a structure of $m^2_{H_u} \sim c\lambda^4 - c'\lambda^2 g^2$ where λ is a coupling constant of Higgs and messenger fields and g is the gauge coupling(s). On the other hand, the two-loop soft mass squared of fields $Q, \bar{U}$ which does not have a direct coupling with messenger has a form of $m^2_{Q_3,\bar{U}_3} \sim -c_3\lambda^2 y_t^2$. This fact was extensively studied in [13]. For sufficiently large λ, large one-loop A terms are generated. At the same time, $m^2_{H_u}$ becomes positive so the soft mass of H_u can be much larger than that in the pure gauge mediation. Moreover, the soft mass of $Q_3, \bar{U}_3$ can be much smaller. If the Higgs H_u superfield directly couples to messengers whereas the top superfields do not,

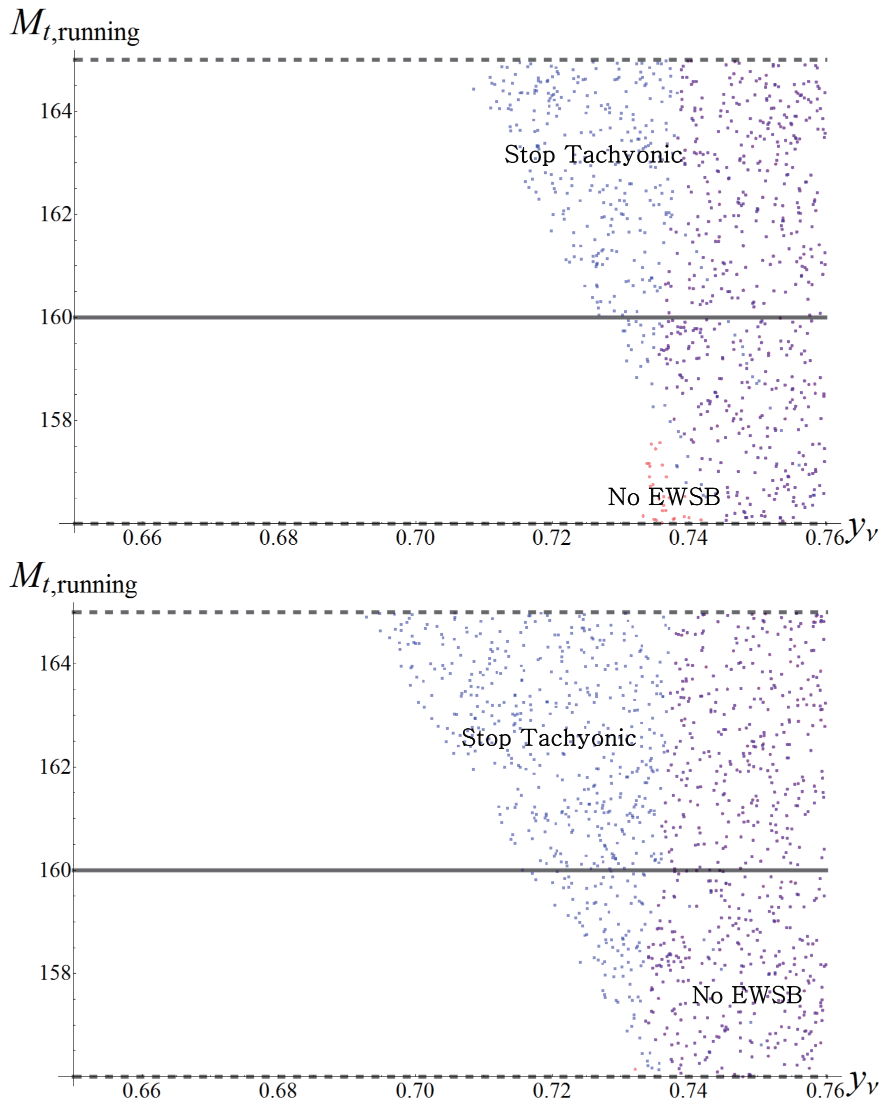

FIG. 2. Phase diagrams indicating stop tachyonic and no EWSB region for $\tan\beta = 10$, $\tan\beta = 30$, respectively. B_N is set to be 5×10^5 GeV.

relatively light stop in natural SUSY can be easily obtained as we can have the small stop soft mass from the effect explained above and the large LR mixing from large A term. A/m^2 problem appears in H_u soft terms such that large A term at the same time generate large $m_{H_u}^2$ at the messenger scale and it can make the electroweak symmetry breaking difficult. It is analogous to the famous $\mu/B\mu$ problem in gauge mediation. To avoid this but to allow the large λ for maximal mixing, large $-c_2\lambda^2 g^2$ part in $m_{H_u}^2$ is required. This can be achieved by introducing new gauge bosons or making strong interaction involved [13]. On the other hand, one loop, negative contribution to $m_{H_u}^2$ can be considered if the messenger scale is low

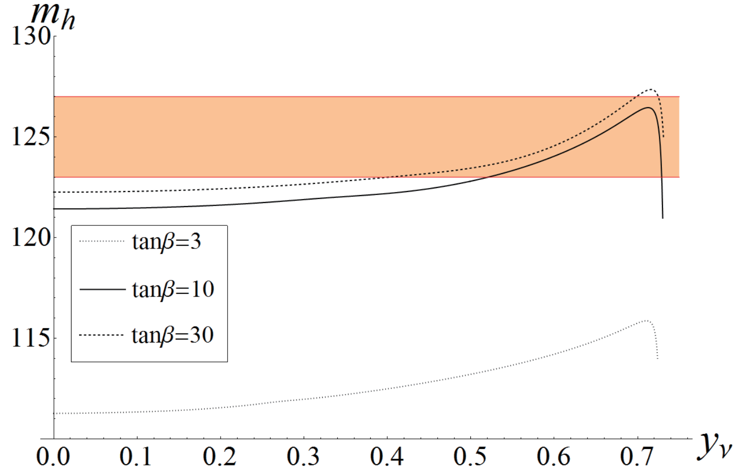

FIG. 3. Higgs mass as a function of y_ν for $B_N = 5 \times 10^5$ GeV, $\rho = 0.1$. Higgs mass can be achieved with the help of Yukawa mediation for large $\tan\beta$ region. At $y_\nu \sim 0.7$, stop mass is approximately 1 TeV.

as analysed in detail in [14].

Neutrino assisted gauge mediation uses the Yukawa coupling among messengers (neutrinos), Higgs and lepton doublets. As a result, Higgs and lepton doublet soft scalar masses get extra contribution from Yukawa mediation. The same A/m^2 problem applies here and neutrino Dirac Yukawa coupling can not be taken to be a large value for successful electroweak symmetry breaking in principle. On the other hand, too large $m_{H_u}^2$ and too large A term may drive stop tachyonic through renormalization group running with top Yukawa. The problem becomes worse as the stop soft scalar mass squared at the messenger scale gets a negative contribution from Yukawa mediation. The situation is shown in Fig. 2. For the running mass of the top quark 160 GeV (the central value), the tachyonic stop appears before the real A/m^2 problem prevents the successful electroweak symmetry breaking as we increase y_ν. The crucial difference compared to the previous work in which A/m^2 problem is emphasized [13, 14] comes from the number of messengers. In neutrino assisted gauge mediation, the number of messengers is three, $N = 3$. The y^2 contribution is effectively replaced by Ny_ν^2. Large N effectively reduces the A/m^2 problem by $1/N$. At the same time smaller y_ν can provide the same impact with the aid of $N > 1$. If tachyonic stop appears as

y_ν gets larger, it would be easy to realise the maximal stop mixing by making the stop soft scalar mass sufficiently small.

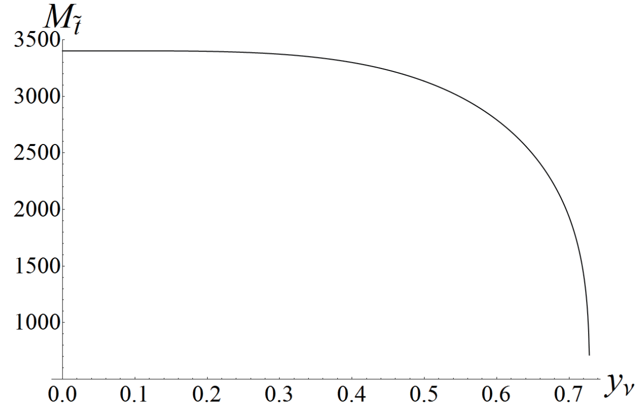

FIG. 4. $M_{\tilde{t}}$ as a function of y_ν for $B_N = 5 \times 10^5$ GeV, $\rho = 0.1$. $\tan\beta = 10$.

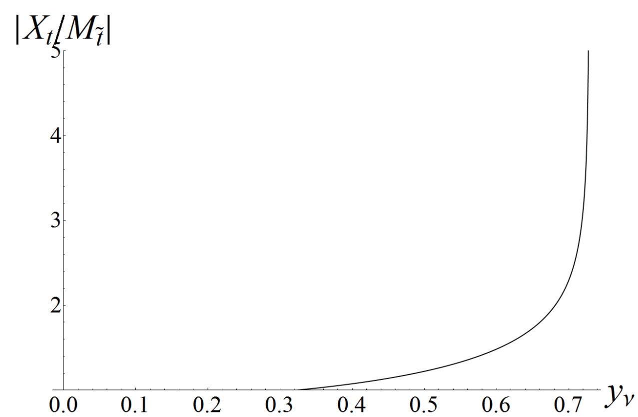

FIG. 5. $X_t/M_{\tilde{t}}$ as a function of y_ν for $B_N = 5 \times 10^5$ GeV, $\rho = 0.1$, $\tan\beta = 10$.

Fig. 3 shows the contribution assisted by neutrino messengers, compared to the minimal gauge mediation which corresponds to $y_\nu = 0$ with stop mass at around 1 TeV. In the minimal gauge mediation, the Higgs mass is computed to be at around $121 \sim 122$ GeV for $\tan\beta = 10 \sim 30$. For $y_\nu = 0.7$, the Higgs mass can be as large as $125 \sim 126$ GeV. 4 to 5 GeV gain in the Higgs mass is obtained in neutrino assisted gauge mediation. The gain does not look impressive but has an impact on allowed superparticle spectrum. In the absence of A_t at the messenger scale as is the case in minimal gauge mediation, this extra 5 to 6 GeV can be achieved by making the logarithmic contribution large and the stop mass should be as heavy as 5 to 10 TeV rather than 2 TeV.

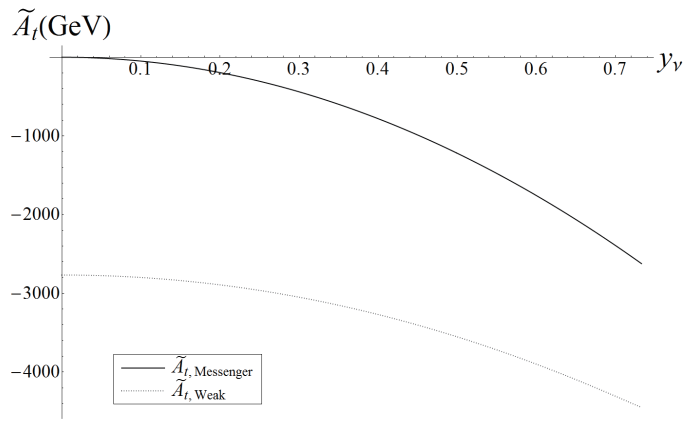

FIG. 6. $\tilde{A}_t \equiv A_t Y_t$ as a function of y_ν for $\tan\beta = 10$, $B_N = 5 \times 10^5$GeV. Without Yukawa mediation, one can obtain $\tilde{A}_t \sim -2700\,$GeV at the weak scale by RG running effects. With help of neutrino mediation at the messenger scale, one can obtain $\tilde{A}_t \sim -4500\,$GeV at weak scale. This drives more stop mixing, which helps 125 GeV Higgs mass.

Note that the plot stops at $y_\nu = 0.72$. Neutrino assisted gauge mediation is classified as Higgs-messenger mixing scenario as the right-handed neutrino is the messenger and the neutrino Dirac Yukawa coupling connects Higgs, lepton doublet and the messenger (right-handed neutrino). The stop soft scalar mass squared gets smaller and becomes tachyonic as the neutrino Dirac Yukawa coupling is increased as in Fig. 4. The logarithmic correction to the Higgs mass also rapidly drops beyond $y_\nu \sim 0.7$ as the stop mass becomes too light (and becomes tachyonic) as is shown in Fig. 3. The maximal mixing is realised around this point, as shown in Fig. 5. This also corresponds to the corner of the parameter space next to the critical point as in [46].

Fig. 6 compares A_t in the minimal gauge mediation and the neutrino assisted gauge mediation both at the messenger scale and the weak scale. Note that A_t by itself is enhanced by 1.5 at the weak scale with the help of messenger scale A_t.

Fig. 7 shows the relation between B_N and the Higgs mass. The neutrino Dirac Yukawa coupling y_ν is chosen to be close to 0.72 which can maximize the Higgs mass for given B_N.

In summary, the minimal gauge mediation needs stop mass at around 5 to 10 TeV to raise the Higgs mass up to 125 GeV. If the right-handed neutrinos are the messengers of the

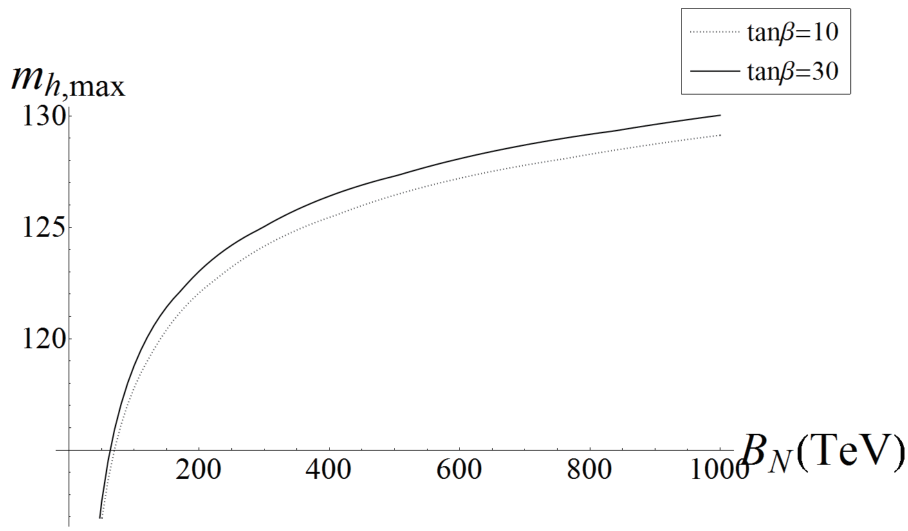

FIG. 7. Maximum values of Higgs mass as a function of B_N. For $\tan\beta = 10$, at least $B_N = 360\,\mathrm{TeV}$ is required to obtain $125\,\mathrm{GeV}$ Higgs mass. For $\tan\beta = 30$, $B_N = 300\,\mathrm{TeV}$ is required. At two points $(\tan\beta = 10, B_N = 360\,\mathrm{TeV})$, $(\tan\beta = 30, B_N = 300\,\mathrm{TeV})$, we display sparticle spectrums in Table III. Also spectrums with $123\,\mathrm{GeV}$ Higgs mass are given for $(\tan\beta = 10, B_N = 240\,\mathrm{TeV})$, $(\tan\beta = 30, B_N = 200\,\mathrm{TeV})$. Sparticle spectrums are displayed in Table IV.

supersymmetry breaking, so called 'neutrino assisted gauge mediation', we can explain 125 GeV Higgs mass with lighter than 2 TeV stop mass.

III. FLAVOR MODEL

In this section, we consider models which can explain neutrino oscillations successfully. Since the SUSY breaking mediation through the neutrino Dirac Yukawa coupling is flavor dependent in general, sizable cLFV could be generated. To avoid this, the neutrino Dirac Yukawa coupling is set to be proportional to the idenitity. In the right-handed neutrino mass basis, it would be proportional to the unitary matrix so soft mass m_L^2, which depends on the combination $Y_\nu^\dagger Y_\nu$ is flavor universal. It is easily achieved by employing the non-abelian discrete symmetry for the tri-bi maximal mixing of the PMNS matrix. Since the tri-bi maximal mixing should be modified to make θ_{13} nonzero, as reported by several

observations[47–51], small corrections should be added. When the neutrino Dirac Yukawa coupling has such corrections, such that it has a deviation from identity, cLFV is generated. We look for several ways to suppress cLFV, at least under the experimental bound.

Superfield	S_4	Z_4	U(1)$_L$	U(1)$_R$
L	$\mathbf{3}$	1	1	1
$\bar{E}$	$\mathbf{2}+\mathbf{1}$	2	-1	0
N	$\mathbf{3}$	3	-1	0
Φ	$\mathbf{3}+\mathbf{3}'$	1	0	0
χ	$\mathbf{1}+\mathbf{2}+\mathbf{3}$	2	2	0
H_u	$\mathbf{1}$	0	0	1
H_d	$\mathbf{1}$	0	0	1
X	$\mathbf{1}$	0	0	2

TABLE I. Charge assignments under $S_4 \times Z_4 \times \mathrm{U}(1)_L \times \mathrm{U}(1)_R$ for leptons, flavons, Higgs, and SUSY breaking spurions.

To make the PMNS matrix tri-bi maximal, we use S_4 discrete symmetry, since it is closely related to the permutation structure of Yukawa couplings. Other discrete symmetries, such as A_4, the even permutation of the S_4 could be used. The main difference is that the first and the second generation of the right-handed leptons belong to $\mathbf{2}$ dimensional representation in S_4 while they correspond to different one dimensional representations, $\mathbf{1}'$, $\mathbf{1}''$ in A_4. In [52–55], the structure we use is obtained from A_4 symmetry and discussion on the deviation from the tri-bi maximal mixing is in parallel. The S_4 symmetry model building is reviewed in [56]. In Appendix A, we summarised representations and tensor products of S_4 group.

For quark sector, the CKM matrix is close to the identity. Deviation from the identity has a hierarchy structure parametrized by some powers of the Cabibbo angle, $\lambda = \sin\theta_C$. On the other hand, the PMNS matrix, mixing matrix in the lepton sector has large mixing angles. Even the smallest mixing angle, θ_{13} is in the order of λ. To explain this, it is natural to assume that $u-$ and $d-$ quark sectors have almost the same structure under the discrete

flavor symmetry whereas the charged lepton and the right-handed neutrino sectors do not. This picture can be realised by introducing appropriate 'flavons' charged under discrete symmetry group and more symmetries can be introduced to forbid useless couplings. Here, we consider the symmetry group $S_4 \times Z_4 \times \mathrm{U}(1)_L$, where $\mathrm{U}(1)_L$ represents a lepton number, which may be discretized. In this paper, we consider superpotential for See-Saw mechanism with flavons Φ and χ,

$$W = -l_{1ij}\bar{E}_i \Phi L_j H_d + l_{2ij} N_i L_j H_u + \frac{1}{2} l_{3ij} X N_i \chi N_j, \tag{19}$$

where $i, j = 1, 2, 3$ are the generation indices and X is a SUSY breaking spurion. For this, S_4, Z_4, $\mathrm{U}(1)_L$ and U(1) R-symmetry quantum numbers are given in Table I.

The charged lepton Yukawa couplings can be constructed from $\bar{E}\Phi L H_d$, the neutrino Dirac Yukawa coupling can be constructed from NL, and the Majorana mass of the heavy neutrinos can be constructed from $XN\chi N$. On the other hand, Φ^2, χ^2, and $\Phi\chi$ cannot couple to the combinations $\bar{E}LH_d$, NL, and XNN to make singlets. Note that $\mathrm{U}(1)_R$ is introduced to forbid unwanted coupling $N\chi N$, which makes B_N in a matrix form, not a constant.

The discrete symmetry quantum number can be extended to the quark sector, such as $Q : (\mathbf{3}, 1, 1, 1, 1)$, $\bar{U} : (\mathbf{2} + \mathbf{1}, 2, 0, 0)$, and $\bar{D} : (\mathbf{2} + \mathbf{1}, 2, 0, 0)$ under $S_4 \times Z_4 \times \mathrm{U}(1)_L \times \mathrm{U}(1)_R$. The flavons $\Phi : (\mathbf{3} + \mathbf{3'}, 1, 0)$ make the singlet combinations $\bar{U}\Phi Q H_u + \bar{D}\Phi Q H_d$ and Yukawa couplings Y_U and Y_D have the same form as the charged lepton Yukawa coupling. They are diagonalized by the same unitary matrix so CKM matrix is the identity in the leading order. If another type of flavon couples to either of up and down quark sectors to give subleading corrections of order λ, it would explain the Cabibbo angle.

Lepton L_i is in the $\mathbf{3}$ and $\bar{E}_j$ is in the $\mathbf{1}$+$\mathbf{2}$ representations, in which $(\bar{E}_1)_{\mathbf{1}} + (\bar{E}_2, \bar{E}_3)_{\mathbf{2}}$. Also there are the SM singlet flavons $\Phi_{\mathbf{3}}$, and $\Phi_{\mathbf{3'}}$ in the $\mathbf{3}$, and $\mathbf{3'}$ representations. We do not provide a complete vacuum alignment in this setup. Instead in Appendix B, we show a few simple examples in which the aligned vacuum is realised. If, for instance, VEVs are arranged to be $\langle \Phi_{\mathbf{3}} \rangle = v_2(1, 1, 1)$, and $\langle \Phi_{\mathbf{3'}} \rangle = v_3(1, 1, 1)$, we have the following Yukawa structure

$$Y_E = \lambda_E \frac{1}{\sqrt{3}} \begin{pmatrix} c & c & c \\ a & a\omega & a\omega^2 \\ b & b\omega^2 & b\omega \end{pmatrix} \tag{20}$$

where $a = (\lambda v_2 + \lambda' v_3)/\Lambda$, $b = (\lambda v_2 - \lambda' v_3)/\Lambda$, $c = \lambda'' v_2/\Lambda$, and $\lambda, \lambda', \lambda''$ are coupling constants of $\bar{E}_{\mathbf{2}} L_{\mathbf{3}} \Phi_{\mathbf{3}}$, $\bar{E}_{\mathbf{2}} L_{\mathbf{3}} \Phi_{\mathbf{3}'}$, and $\bar{E}_{\mathbf{1}} L_{\mathbf{3}} \Phi_{\mathbf{3}}$, respectively. In this case, $Y_E^\dagger Y_E$ has the form of

$$Y_E^\dagger Y_E = |\lambda_E|^2 \begin{pmatrix} a^2+b^2+c^2 & c^2+a^2\omega+b^2\omega^2 & c^2+b^2\omega+a^2\omega^2 \\ c^2+a^2\omega^2+b^2\omega & a^2+b^2+c^2 & c^2+a^2\omega+b^2\omega^2 \\ c^2+b^2\omega^2+a^2\omega & c^2+a^2\omega^2+b^2\omega & a^2+b^2+c^2 \end{pmatrix}, \tag{21}$$

which will be diagonalized to $|\lambda_E|^2((\epsilon^3)^2, (\epsilon)^2, 1)$ by the unitary matrix,

$$V_L^l = \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 1 & 1 \\ 1 & \omega^2 & \omega \\ 1 & \omega & \omega^2 \end{pmatrix}. \tag{22}$$

Here we use $\epsilon \simeq m_\mu/m_\tau$ as the order parameter. Then, $c = \epsilon^3$, $a = \epsilon$ and $b = 1$.

On the other hand, let heavy neutrinos N_i be in the triplet $\mathbf{3}$. $\Phi_{\mathbf{3}}$ and $\Phi_{\mathbf{3}'}$ cannot couple to the combination $L_i N_j$ by Z_4 and $\mathrm{U}(1)_L$ symmetries as well as the SM gauge symmetry. Since the combination $L_1 N_1 + L_2 N_2 + L_3 N_3$ is a singlet, we naturally have the neutrino Dirac Yukawa coupling Y_ν proportional to the identity. Finally, $X N_i N_j$ has again the form of $\mathbf{3} + \mathbf{3}' + \mathbf{1} + \mathbf{2}$. Φ's cannot couple to it while singlet $\chi_{\mathbf{1}}$ and triplet $\chi_{\mathbf{3}}$ in the singlet and triplet representation can do, so we have the following Majorana mass term:

$$M_N = \begin{pmatrix} w_1 & 0 & w_2 \\ 0 & w_1 & 0 \\ w_2 & 0 & w_1 \end{pmatrix} \tag{23}$$

where $\langle \chi_{\mathbf{1}} \rangle = w_1$ and $\langle \chi_{\mathbf{3}} \rangle = w_2(0,1,0)$, respectively. Therefore, the neutrino mass matrix $M_\nu = -v_u^2 Y_\nu^T M_N^{-1} Y_\nu$ is diagonalized by

$$V_L^\nu = \begin{pmatrix} \frac{1}{\sqrt{2}} & 0 & -\frac{1}{\sqrt{2}} \\ 0 & 1 & 0 \\ \frac{1}{\sqrt{2}} & 0 & \frac{1}{\sqrt{2}} \end{pmatrix} \tag{24}$$

so we obtain the PMNS matrix in the tri-bi maximal mixing,

$$V_{\mathrm{PMNS}} \equiv (V_L^l)^\dagger V_L^\nu = \begin{pmatrix} \sqrt{\frac{2}{3}} & \frac{1}{\sqrt{3}} & 0 \\ -\omega \frac{1}{\sqrt{6}} & \omega \frac{1}{\sqrt{3}} & e^{-i5\pi/6} \frac{1}{\sqrt{2}} \\ -\omega^2 \frac{1}{\sqrt{6}} & \omega^2 \frac{1}{\sqrt{3}} & e^{i5\pi/6} \frac{1}{\sqrt{2}} \end{pmatrix}. \tag{25}$$

In this construction, S_4 triplet flavons have VEVs in the direction of $(1,1,1)$ or $(0,1,0)$. These directions are easily stabilized compared to other directions, such as $(1,1,0)$, as argued in Appendix B. Note that Y_ν proportional to the identity does not give rise to LFV. In Eq. (12), we see m_L^2 from the neutrino Dirac Yukawa mediation is flavor universal. In the right-handed neutrino and the charged lepton mass basis, Y_ν moves to $V_L^\nu Y_\nu$ and m_L^2 moves to $(V_L^l)^\dagger m_L^2 V_L^l$. As a result, PMNS matrix is mulitplied and will change m_L^2 matrix. However, if the neutrino Dirac Yukawa matrix is proportional to the identity matrix, the property of $V\mathbb{I}V^\dagger = \mathbb{I}VV^\dagger = \mathbb{I}$ cancels out such effects.

There are various ways to put corrections to make non-zero θ_{13}. Moreover, corrected neutrino mass matrix should be consistent with the measurements of θ_{12}, θ_{23} as well as neutrino mass squared differences, $\Delta m^2_{\rm sol} \equiv m_2^2 - m_1^2$ and $|\Delta m^2_{\rm atm}| \equiv |m_3^2 - m_2^2|$. Since the overall neutrino mass scale is not known, the important quantity is the ratio of neutrino mass squared differences, as described in [55],

$$\sqrt{|R|} \equiv \sqrt{\frac{\Delta m^2_{\rm atm}}{\Delta m^2_{\rm sol}}}. \tag{26}$$

The measured values adopted in [57] are given by

$$\begin{aligned} \Delta m^2_{\rm sol} &= (7.50 \pm 0.20) \times 10^{-5} {\rm eV}^2 \\ \Delta m^2_{\rm atm} &= (0.00232)^{+0.00012}_{-0.00008} {\rm eV}^2 \end{aligned}$$

$$\begin{aligned} \sin^2(2\theta_{12}) &= 0.857 \pm 0.024 \\ \sin^2(2\theta_{23}) &> 0.95 \\ \sin^2(2\theta_{13}) &= 0.098 \pm 0.013 \end{aligned} \tag{27}$$

in the 90% C. L. The global analysis for such quantities can be found in [58, 59].

Suppose, for simplicity, we leave the charged lepton sector untouched and correct neutrino sector only. Moreover, we keep the mixings of ν_2 with $\nu_{1,3}$ forbidden, so that V_L^ν is modified to

$$V_L^\nu = \begin{pmatrix} \cos(\frac{\pi}{4}+\delta) & 0 & -\sin(\frac{\pi}{4}+\delta) \\ 0 & 1 & 0 \\ \sin(\frac{\pi}{4}+\delta) & 0 & \cos(\frac{\pi}{4}+\delta) \end{pmatrix}. \tag{28}$$

For small δ, $\cos(\frac{\pi}{4}+\delta) \simeq (1/\sqrt{2})(1-\delta)$ and $\sin(\frac{\pi}{4}+\delta) \simeq (1/\sqrt{2})(1+\delta)$. From

$$\begin{aligned} V_{\rm PMNS} &= (V_L^l)^\dagger V_L^\nu \\ &= \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 1 & 1 \\ 1 & \omega & \omega^2 \\ 1 & \omega^2 & \omega \end{pmatrix} \frac{1}{\sqrt{2}} \begin{pmatrix} (1-\delta) & 0 & -(1+\delta) \\ 0 & 1 & 0 \\ (1+\delta) & 0 & (1-\delta) \end{pmatrix}, \end{aligned} \tag{29}$$

we see (13) element of the PMNS matrix is given by

$$|V_{e3}| = \left|\frac{2\delta}{\sqrt{6}}\right|. \tag{30}$$

If such corrections are entirely present in the right-handed neutrino Majorana mass term while Y_ν is untouched, there would be no observable charged lepton flavor violating process. For example, let us introduce a doublet flavon $\chi_{\mathbf{2}}$. Then, its VEV modifies the diagonal elements of the Majorana mass matrix. With $\langle\chi_2\rangle = x^2(1,1)$, diagonal term has a correction $x^2[2N_1N_1 - N_2N_2 - N_3N_3]$. In principle, by introducing several doublets with different VEVs, each diagonal term can be different.

A. Model I

Besides putting correction to M_N, one can find S_4 doublet VEVs giving corrections to Y_ν to make a sizable θ_{13} while dangerous charged lepton flavor violation is suppressed. To see this, consider the general S_4 doublet VEV, (a,b) where a and b are complex numbers. With this VEV and coupling λ_1, Y_ν can be modified as

$$\begin{pmatrix} 1+\lambda_1(a+b) & 0 & 0 \\ 0 & 1+\lambda_1(b\omega+a\omega^2) & 0 \\ 0 & 0 & 1+\lambda_1(b\omega^2+a\omega) \end{pmatrix} \tag{31}$$

In this case, $Y_\nu^\dagger Y_\nu$ in the charged lepton mass basis is given by

$$\begin{pmatrix} 1+\lambda_1^2(|a|^2+|b|^2) & \lambda_1(a^*+b)+\lambda_1^2ab^* & \lambda_1(a+b^*)+\lambda_1^2a^*b \\ \lambda_1(a+b^*)+\lambda_1^2a^*b & 1+\lambda_1^2(|a|^2+|b|^2) & \lambda_1(a^*+b)+\lambda_1^2ab^* \\ \lambda_1(a^*+b)+\lambda_1^2ab^* & \lambda_1(a+b^*)+\lambda_1^2a^*b & 1+\lambda_1^2(|a|^2+|b|^2) \end{pmatrix}. \tag{32}$$

If $\lambda_1(a^*+b)+\lambda_1^2ab^* = 0$, all the off diagonal elements vanish. For example, $\lambda_1 = 1$ and $a = b = \omega$ is the case. This condition also implies that off diagonal terms of $Y_\nu^\dagger Y_\nu Y_\nu^\dagger Y_\nu$ vanish

so we do not expect any sizeable cLFV. However, this condition requires a cancellation of two different flavon contributions and is considered as a serious fine tuning different from vacuum alignment. We do not pursue this possibility any longer in this paper.

If $a^* = -b$ and both $|a|$ and $|b|$ are smaller than one, the (12) element of Y_ν is given by $-\lambda_1^2(a^*)^2$. The (23) element is the same and the (13) element is its complex conjugate, $-\lambda_1^2 a^2$. In this way, LFV is suppressed quadratically even though it does not vanish. For $Y_\nu^\dagger Y_\nu Y_\nu^\dagger Y_\nu$ term, the (12) element is $2[1+\lambda_1^2(|a|^2+|b|^2)][\lambda_1(a^*+b)+\lambda_1^2 ab^*]+[\lambda_1(a+b^*)+\lambda_1^2 a^* b]^2$. For $Y_\nu^\dagger Y_\nu$ term, the (23) element is the same and the (13) element is its complex conjugate. When $a^* = -b$, it is $-2\lambda_1^2 a^2(1+2\lambda_1^2|a|^2)+\lambda_1^4(a^*)^4$, which is quadratically suppressed for small a. For illustration, suppose $\lambda_1 a = \lambda_1 b = i\rho$. The stabilization of such doublet VEV is discussed in Appendix B. The neutrino Dirac Yukawa has the form of

$$Y_\nu = y_\nu \begin{pmatrix} 1+2i\rho & 0 & 0 \\ 0 & 1-i\rho & 0 \\ 0 & 0 & 1-i\rho \end{pmatrix} \tag{33}$$

and off-diagonal terms of $Y_\nu^\dagger Y_\nu$ in the charged lepton mass basis is suppressed to $\mathcal{O}(\rho^2)$, as expected,

$$(V_L^l)^\dagger (Y_\nu^\dagger Y_\nu) V_L^l = |y_\nu|^2 \begin{pmatrix} 1+2\rho^2 & \rho^2 & \rho^2 \\ \rho^2 & 1+2\rho^2 & \rho^2 \\ \rho^2 & \rho^2 & 1+2\rho^2 \end{pmatrix}. \tag{34}$$

With this Y_ν, neutrino mass matrix is given by

$$M_\nu = -|y_\nu|^2 \frac{v^2 \sin^2\beta}{2w_1} \frac{1}{1-x^2} \begin{pmatrix} 1+4i\rho-4\rho^2 & 0 & -x(1+i\rho+2\rho^2) \\ 0 & (1-x^2)(1-2i\rho-\rho^2) & 0 \\ -x(1+i\rho+2\rho^2) & 0 & 1-2i\rho-\rho^2 \end{pmatrix} \tag{35}$$

and the deviation of mixing from $\pi/4$ is given by

$$\delta = \left| \frac{-6i\rho+3\rho^2}{4x(1+i\rho+2\rho^2)} \right| \simeq \frac{3\rho}{2x} \tag{36}$$

such that

$$|V_{e3}| \simeq \frac{3\rho}{\sqrt{6}x}. \tag{37}$$

To the first order in ρ, mass eigenvalues are given by

$$-|y_\nu|^2\frac{v^2\sin^2\beta}{2w_1}\Big(\frac{1+i\rho}{1+x},1-2i\rho,\frac{1+i\rho}{1-x}\Big). \tag{38}$$

Taking absolute values of these eigenvalues, we obtain neutrino masses $-[|y_\nu|^2v^2\sin^2\beta/(2w_1)](1/(1+x),1,1/(1-x))+\mathcal{O}(\rho^2)$.

In summary, we expect that even though the charged lepton flavor violating effects are generated in the A_E term at one loop and in the m_L^2 term at two loop, they can be suppressed by extra small expansion parameter ρ proportional to θ_{13}. With the vacuum alignment of the doublet flavon $i(v,v)$, it is possible to cancel the first order correction of ρ and the off-diagonal elements of the slepton mass squared would have ρ^2 suppression as a result. Fig. 8 shows how measured θ_{13} can be explained for the choices of ρ and x parameters satisfying observed neutrino mass squared ratio, $\sqrt{R}$. The observed $\theta_{13}\sim 0.15$ can be accommodated for $\rho\sim 0.1$.

B. Model II

Of course, θ_{13} can come from both Majorana mass correction and neutrino Dirac Yukawa correction. Only the neutrino Dirac Yukawa coupling can affect the cLFV. To see the two-parameter case, consider

$$Y_\nu = y_\nu\begin{pmatrix} 1+2i\rho & 0 & 0\\ 0 & 1-i\rho & 0\\ 0 & 0 & 1-i\rho\end{pmatrix} \tag{39}$$

and

$$M_N=\begin{pmatrix} w_1 & 0 & w_2\\ 0 & w_1 & 0\\ w_2 & 0 & w_1(1-\zeta)\end{pmatrix}. \tag{40}$$

The neutrino mass is given by

$$M_\nu=-|y_\nu|^2\frac{v^2\sin^2\beta}{2w_1}\frac{1}{1-x^2-\zeta}\begin{pmatrix} 1+4i\rho-\zeta & 0 & -x(1+i\rho)\\ 0 & (1-x^2)(1-2i\rho)-\zeta & 0\\ -x(1+i\rho) & 0 & 1-2i\rho\end{pmatrix}+\mathcal{O}(\rho^2,\zeta^2) \tag{41}$$

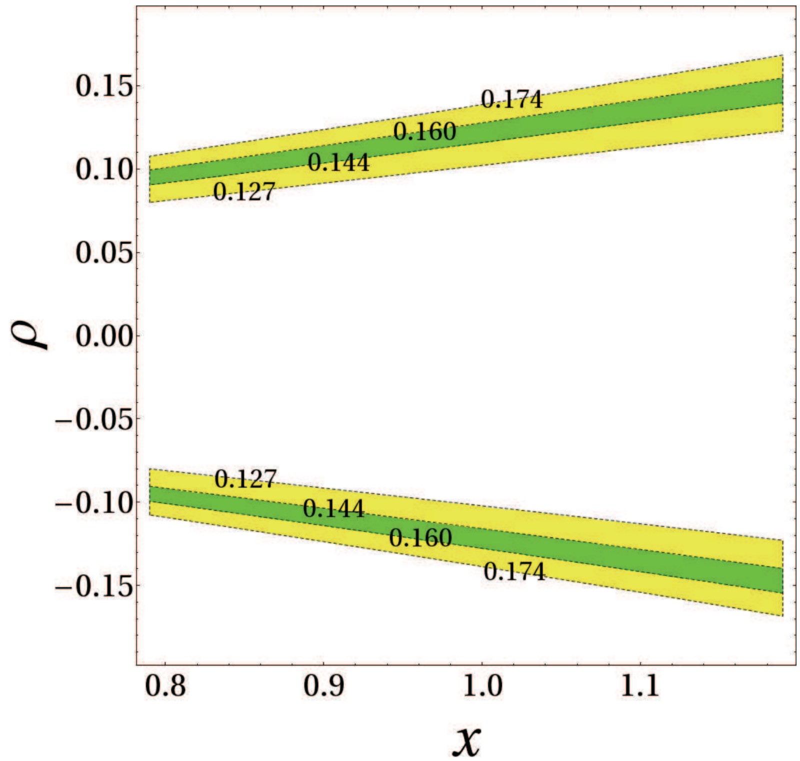

FIG. 8. θ_{13} with respect to ρ and x parameters. All points in the colored region satisfy neutrino oscillation experiments. Neutrino θ_{13}, indicated on contour label in radian, is measured as $0.144 < \theta_{13} < 0.160$ in 1σ level, $0.127 < \theta_{13} < 0.174$ in 3σ level.

where $x = w_2/w_1$ again. Then, three neutrino mass eigenvalues are given by

$$-|y_\nu|^2 \frac{v^2 \sin^2\beta}{2w_1}\Big(\frac{1+i\rho}{1+x}+\frac{\zeta}{2(1+x)^2}, 1-2i\rho, \frac{1+i\rho}{1-x}+\frac{\zeta}{2(1-x)^2}\Big) \tag{42}$$

and

$$\delta = \frac{\sqrt{36\rho^2+\zeta^2}}{4x}. \tag{43}$$

Hence, we see the (13) element of the PMNS matrix is given by

$$|V_{e3}| = \Big|\frac{2\delta}{\sqrt{6}}\Big| = \Big|\frac{\sqrt{36\rho^2+\zeta^2}}{2\sqrt{6}x}\Big|. \tag{44}$$

Moreover, m_L^2 from Yukawa mediation is controlled by the parameter ρ only and $(V_L^l)^\dagger(Y_\nu^\dagger Y_\nu)V_L^l$ is the same as the previous case,

$$(V_L^l)^\dagger(Y_\nu^\dagger Y_\nu)V_L^l = |y_\nu|^2 \begin{pmatrix} 1+2\rho^2 & \rho^2 & \rho^2 \\ \rho^2 & 1+2\rho^2 & \rho^2 \\ \rho^2 & \rho^2 & 1+2\rho^2 \end{pmatrix}. \tag{45}$$

In the limit of $\zeta \to 0$, both θ_{13} and cLFV come from the neutrino Dirac Yukawa which corresponds to the Model I. In the opposite limit, $\rho \to 0$, θ_{13} is entirely obtained from Majorana mass term and cLFV does not appear.

In addition, we can also constrain absolute mass scale of light neutrinos. The most stringent constraint on neutrino absolute mass is given by CMB data of WMAP experiment, combined with supernovae data and data on galaxy clustering, $\Sigma_j m_j \lesssim 0.68\text{eV}$, 95% C.L. Conservatively, we set the bound $2.6 \times 10^{14}\,\text{GeV} \lesssim M_N$. Throughout paper, we use $M_N = 5 \times 10^{14}\text{GeV}$, the heaviest right-handed neutrino mass.

IV. CHARGED LEPTON FLAVOR VIOLATION

Since flavor structures of supersymmetric particles can be different from those of SM partners, flavor number is easily violated in SUSY. In general, the structure of the slepton mass matrix raises dangerous cLFV. Such cLFV in SUSY is studied in [60, 61]. In our model, once the identity structure of the neutrino Yukawa coupling Y_ν is broken, cLFV is produced. As a possible modification, one may put off-diagonal terms into Y_ν. On the other hand, when the degeneracy of Y_ν is broken, the combination of Y_νs in the charged lepton mass basis, $(V_L^l)^\dagger (Y_\nu^\dagger Y_\nu) V_L^l$ has off-diagonal terms as shown in Sec. III. The slepton mass squared gets extra contribution from neutrino Dirac Yukawa interactions,

$$\delta m_L^2 = \frac{B_N^2}{(4\pi)^4}\Big[\Big(\text{Tr}[Y_\nu Y_\nu^\dagger] + 3\text{Tr}[Y_U Y_U^\dagger] - 3g_2^2 - \frac{1}{5}g_1^2\Big) Y_\nu^\dagger Y_\nu + 3Y_\nu^\dagger Y_\nu Y_\nu^\dagger Y_\nu\Big]. \tag{46}$$

In the charged lepton mass basis, $(V_L^l)^\dagger m_L^2 V_L^l$ has off-diagonal elements and cLFV appears. Even though this is a general feature, it is also possible to find some parameter space in which charged lepton number is conserved. For example, in Sec. III A, off diagonal terms of the slepton soft mass squared, $(m_L^2)_{12}$ can vanish for specific value of y_ν. Corresponding condition would be

$$(\delta m_L^2)_{12} \propto \Big[\text{Tr}[Y_\nu^\dagger Y_\nu] + 3\text{Tr}[Y_U Y_U^\dagger] - 3g_2^2 - \frac{1}{5}g_1^2\Big](Y_\nu^\dagger Y_\nu)_{12} + 3(Y_\nu^\dagger Y_\nu Y_\nu^\dagger Y_\nu)_{12} = 0, \tag{47}$$

which is equivalent to

$$\begin{aligned}(\delta m_L^2)_{12} &\propto \Big[3(1+2\rho^2)y_\nu^2 + 3y_t^2 - 3g_2^2 - \frac{1}{5}g_1^2\Big] y_\nu^2\rho^2 + 3y_\nu^4 2\rho^2 \\ &\simeq y_\nu^2\Big[3y_t^2 - 3g_2^2 - \frac{1}{5}g_1^2\Big]\rho^2 + 9y_\nu^4\rho^2 + \mathcal{O}(\rho^4) = 0.\end{aligned} \tag{48}$$

Near the GUT scale, $g_1^2 \simeq g_2^2 \simeq 4\pi/28$, and $y_t \simeq 0.5$ so off diagonal term vanishes for $y_\nu \simeq 0.28$. For this value of Y_ν, there would be no unwanted cLFV. This is different from the condition that diagonal contribution involving Y_ν vanishes,

$$\begin{aligned}(\delta m_L^2)_{ii} &\propto \Big[3(1+2\rho^2)y_\nu^2 + 3y_t^2 - 3g_2^2 - \frac{1}{5}g_1^2\Big]y_\nu^2(1+2\rho^2) + 3y_\nu^4(1+8\rho^2) \\ &\simeq y_\nu^2\Big[3y_t^2 - 3g_2^2 - \frac{1}{5}g_1^2\Big] + 6y_\nu^4 + \mathcal{O}(\rho^2) = 0\end{aligned} \tag{49}$$

which is satisfied for $y_\nu \simeq 0.34$.

Of course, it does not mean that y_ν should take the lepton number conserving value. We have many constraints on y_ν from various observations. In this paper, we try to explain the 125GeV Higgs mass with large A term generated from y_ν. On the other hand, one may try to explain deviation of muon $g-2$ from the SM prediction. Moreover, degeneracy breaking parameter ρ is used to explain sizable θ_{13}. However, it is also difficult to find an appropriate value of y_ν which satisfies all of them. In this section, we present the cLFVs for parameters explaining the 125GeV Higgs mass with large A term and θ_{13}. Thereafter, we visit the muon $g-2$ constraints and the relation among θ_{13}, cLFV, and the Higgs mass.

A. Experimental status

The current experimental bounds and future sensitivities for various cLFV processes in the 90% C. L. are summarised in Table II[57, 62].

B. $l_j \to l_i\gamma$

The amplitude for $l_j \to l_i\gamma$ is written as

$$T = e\epsilon^{\mu*}\overline{u_i}(p-q)\Big[q^2\gamma_\mu(A_1^L P_L + A_1^R P_R) + m_{l_j} i\sigma_{\mu\nu}q^\nu(A_2^L P_L + A_2^R P_R)\Big]u_j(p). \tag{50}$$

On the mass shell ($q^2 \to 0$), gauge invariance imposes that the chirality preserving part does not contribute to the $l_j \to l_i\gamma$ process. Hence, chirality flipping should take place in the on-shell $l_j \to l_i\gamma$ process. The decay rate is given by

$$\Gamma(l_j \to l_i\gamma) = \frac{e^2}{16\pi}m_{l_j}^5(|A_2^L|^2 + |A_2^R|^2) \tag{51}$$

Observables	Experimental bound	Future sensitivity
$\mathrm{Br}(\mu \to e\gamma)$	2.4×10^{-12}[63]	$\mathcal{O}(10^{-13})$ [63]
$\mathrm{Br}(\tau \to \mu\gamma)$	4.4×10^{-8}[64]	2.4×10^{-9}[69]
$\mathrm{Br}(\tau \to e\gamma)$	3.3×10^{-8}[64]	3.0×10^{-9} [69]
$\mathrm{Br}(\mu \to 3e)$	1.0×10^{-12} [65]	$\mathcal{O}(10^{-16})$ [70]
$\mathrm{Br}(\tau \to 3e)$	2.7×10^{-8}[66]	2.3×10^{-10} [69]
$\mathrm{Br}(\tau \to 3\mu)$	2.1×10^{-8}[66]	8.2×10^{-10}[69]
$\frac{\Gamma(\mu\mathrm{Ti}\to e\mathrm{Ti})}{\Gamma(\mu\mathrm{Ti}\to\mathrm{capture})}$	4.3×10^{-12}[67]	$\mathcal{O}(10^{-18})$[71]
$\frac{\Gamma(\mu\mathrm{Au}\to e\mathrm{Au})}{\Gamma(\mu\mathrm{Au}\to\mathrm{capture})}$	7.0×10^{-13}[68]	

TABLE II. Various LFV experimental bounds and future sensitivities. The table is adopted from [72].

and the branching ratio yields approximately

$$\mathrm{Br}(l_j \to l_i\gamma) \sim \frac{\alpha^3}{G_F^2}\frac{1}{m_{\mathrm{SUSY}}^4}\Big(\frac{(m_L^2)_{ij}}{m_{\mathrm{SUSY}}^2}\Big)^2. \tag{52}$$

In the mass insertion scheme, the chirality flipping can be easily analyzed. Consider first the case of the neutralino-charged slepton internal loop, as shown in Fig. 9. Fig. 9 (a) shows the chirality flipping from a fermion mass insertion in the external lepton line. In Fig. 9 (b), chirality flipping takes place in the slepton internal line through the LR mixing insertion, $m_j(A - \mu\tan\beta)$. This term consists of flavor universal part $-m_j\mu\tan\beta$, which can be enhanced in the limit of large $\tan\beta$ and large μ. The chirality flipping in Fig. 9 (c) is given by the Yukawa coupling of the lepton-slepton-Higgsino vertex. This vertex contains $1/\cos\beta$ factor which combines with a $\sin\beta$ in the Higgsino-gaugino mixing to give a $\tan\beta$ dependence to the diagram. Therefore, this diagram is enhanced in the large $\tan\beta$ limit. Note that it is inversely proportional to the μ, the Higgsino mass. Since this diagram contains SUSY mass scale only, unlike other diagrams proportional to the Higgs VEV v through m_j, it is dominant over all other diagrams with the neutralino-charged slepton internal loop in many cases. However, since the Higgsino-bino mass insertion $M_Z\sin\beta\sin\theta_W$ and Higgsino-

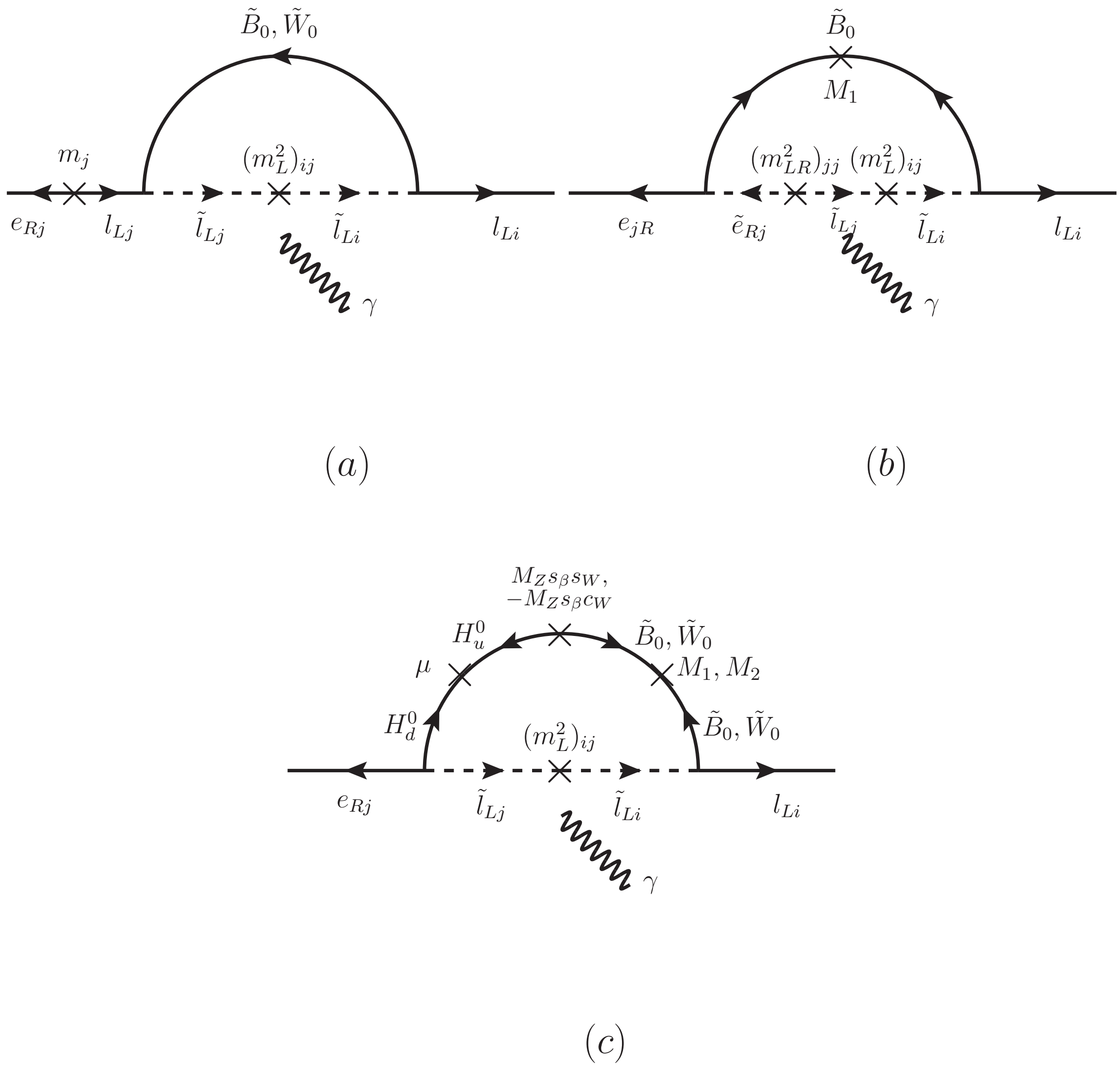

FIG. 9. Feynman diagrams for $l_j \to l_i\gamma$ process with neutralino-charged slepton internal lines in the mass insertion scheme.

wino mass insertion $-M_Z \sin\beta\cos\theta_W$ have the opposite signs, slight destructive interference occurs.

Next, the case of the chargino-sneutrino internal loop is shown in Fig. 10. Diagrams are similar to those of the neutralino-charged slepton internal loop, except the absence of the slepton LR mixing, since the right handed neutrinos are already integrated out. Chirality flipping can occur either in the external lepton line (Fig. 10 (a)) or in the lepton-sneutrino-Higgsino vertex(Fig. 10 (b)). The latter diagram dominates over the former one, and since it does not have a destructive interference, it becomes the leading contribution over all other

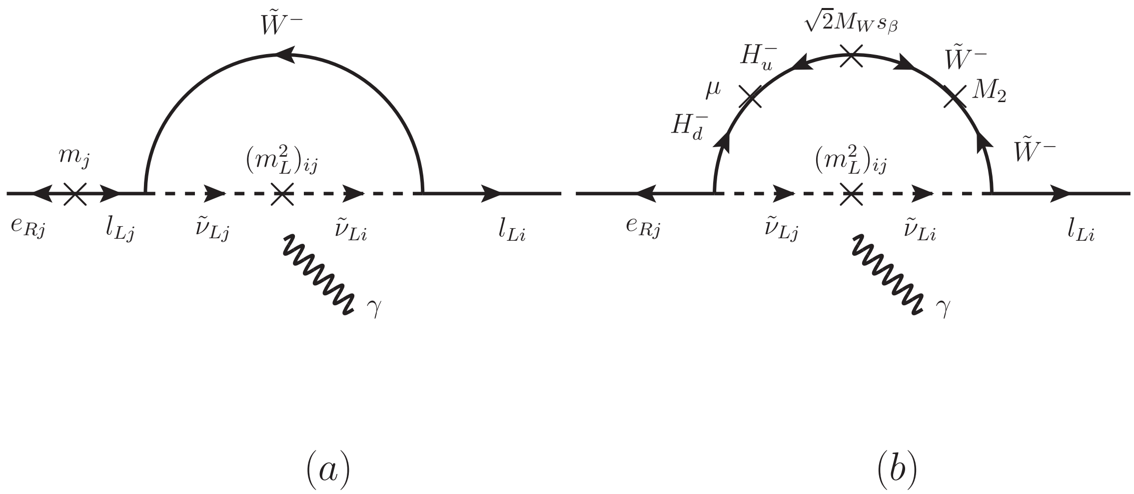

FIG. 10. Feynman diagrams for $l_j \to l_i\gamma$ process with chargino-sneutrino internal lines in the mass insertion scheme.

diagrams in many cases. The similar argument also applies to the discussion of muon $g-2$, whose SUSY contribution comes from the same diagram with flavor conservation. Following this diagram, SUSY enhances the muon $g-2$ for positive μ[73, 74].

In Fig. 11, we show branching ratios of various $l_j \to l_i\gamma$ processes for Sec. III A. In the graph, neutrino Dirac Yukawa couplings are fixed to be $y_\nu = 0.65$ and $\rho = 0.1$, while $\tan\beta$ and SUSY breaking scale are varied. Since off-diagonal terms of m_L^2 in the charged lepton mass basis are the same, normalized branching ratio, $\Gamma(l_j \to l_i\gamma)/m_j^5$ are almost identical. Therefore, branching ratios are closely related to the total decay rate of mother particle. For example, since total decay rate of tau is about 5.3 times larger than that of muon, branching ratio of $\mathrm{Br}(\mu \to e\gamma)$ is about 5.3 times larger than $\mathrm{Br}(\tau \to e\gamma)$ and $\mathrm{Br}(\tau \to \mu\gamma)$ which are almost the same.

C. $l_j^- \to l_i^- l_i^- l_i^+$

In many cases, dominant contribution comes from the photon penguin. Z boson penguin is suppressed in general because of the accidental cancellation when the neutralino or chargino is pure gaugino or pure Higgsino[75]. Such accidental cancellation is broken by introducing TeV scale physics which couples to the sneutrino with the sizable coupling. This can be realised in the R-parity violating model or in the TeV inverse seesaw, for

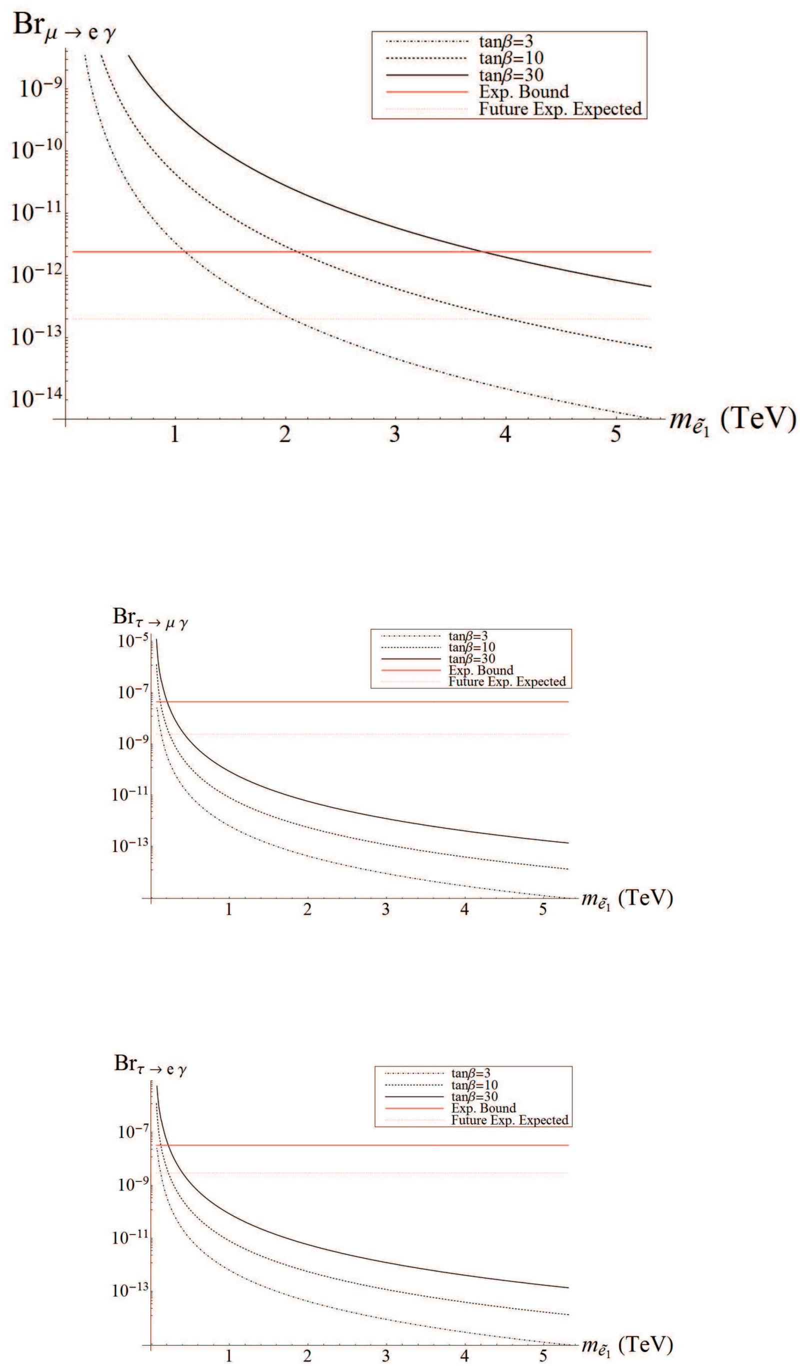

FIG. 11. Branching ratios of $\mu \to e\gamma$, $\tau \to \mu\gamma$ and $\tau \to e\gamma$ with respect to the lightest selectron mass for $\tan\beta = 3, 10, 30$, $y_\nu = 0.65$ and $\rho = 0.1$.

example[72, 75].

In our case, photon penguin is a leading contribution, so we have a simple relation between $\mathrm{Br}(l_j \to l_i\gamma)$,

$$\frac{\mathrm{Br}(l_j \to 3l_i)}{\mathrm{Br}(l_j \to l_i\gamma)} = \frac{\alpha}{3\pi}\Big(\ln\frac{m_{l_j}^2}{m_{l_i}^2} - \frac{11}{4}\Big). \tag{53}$$

The box diagram is suppressed in general, except for some special cases, such as in SUSY with Dirac gauginos[76].

In Fig. 12, we show branching ratios of various $l_j \to 3l_i$ processes for for Sec. III A. Fixed parameters are the same as $l_j \to l_i\gamma$ process. We see that $\mathrm{Br}(\mu \to 3e)$ is about 0.018 times suppressed than $\mathrm{Br}(\mu \to e\gamma)$ so Eq. (53) is satisfied. Photon penguin is a leading contribution for $\mu \to 3e$ process. In the absence of special characteristic which can overcome the natural size of the branching ratio, $\mathrm{Br}(l_j^- \to l_i^- l_i^- l_i^+)$ is α/π suppressed compared to $\mathrm{Br}(l_j^- \to l_i^-\gamma)$.

D. $\mu - e$ conversion

Conversion of the stopped muons in a nuclei to the electron is a promising channel to look for the charged lepton flavor violation. In principle there are many different operators including scalar, photon mediated vector, Z-boson mediated vector operators in addition to the dipole operator. Muon to electron conversion rate is conventionally normalised by muon capture rate.

$$B_{\mu\to e}(Z) = \frac{\Gamma_{\text{conv}}(Z,A)}{\Gamma_{\text{capt}}(Z,A)}. \tag{54}$$

Here Z is the atomic number of the atom. Different target provide a different $B_{\mu\to e}(Z)$ and relative ratio of $B_{\mu\to e}(Z)$ of at least two different target can provide information on possible types of the operators as different operators predict different ratios. In supersymmetric models[77], dominant contribution to $\mu - e$ conversion comes from the dipole operator. As a result, $B(\mu \to e)(Z)$ is predicted to be suppressed by α/π compared to $B(\mu \to e\gamma)$. For different choice of Z, the conversion is suppressed by $10^{-3} \sim 5 \times 10^{-3}$. Current limit on the conversion rate is comparable to $\mu \to e\gamma$ branching ratio, but the future experiments on μ to e conversion will have better sensitivity. We plot $\mu - e$ conversion rate with the expected future sensitivity of planned experiments in Fig.13.

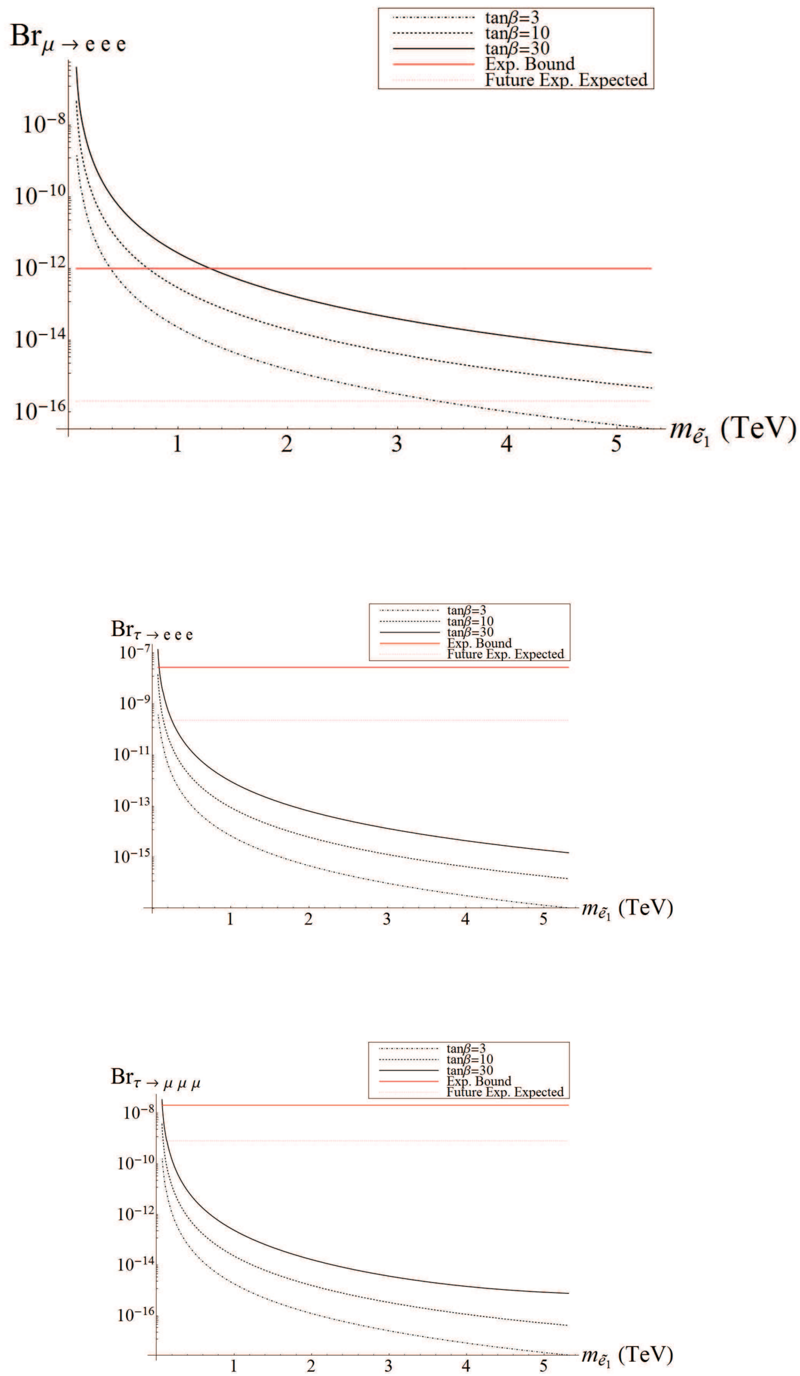

FIG. 12. Branching Ratios of $\mu \to 3e$, $\tau \to 3e$ and $\tau \to 3\mu$ with respect to the lightest selectron mass for $\tan\beta = 3, 10, 30$, $y_\nu = 0.65$ and $\rho = 0.1$.

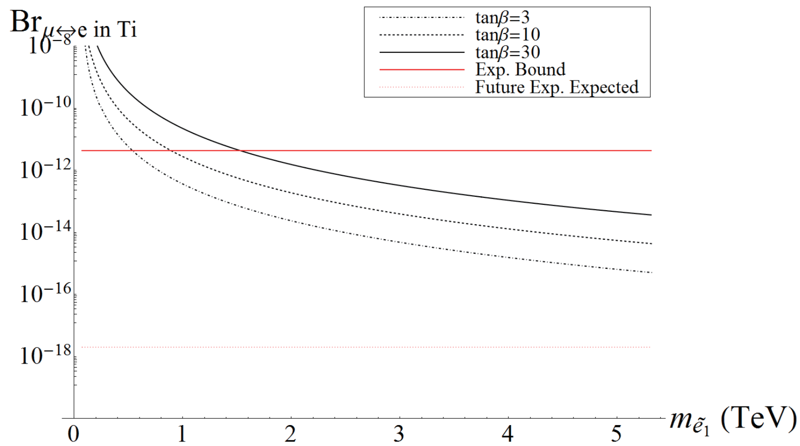

FIG. 13. $\mu - e$ conversion rate with respect to the lightest selectron mass for $\tan\beta = 3, 10, 30$, $y_\nu = 0.65$ and $\rho = 0.1$.

E. Correlation between Muon $g-2$, θ_{13}, cLFV and the Higgs

The anomalous magnetic moment of muon (muon $g-2$) has a long standing sizable deviation from the SM prediction. The observed value is [78]

$$a_\mu(\text{Exp}) = 11659208.9(6.3) \times 10^{-10}, \tag{55}$$

whereas the SM prediction[79] is given by

$$a_\mu(\text{SM}) = 11659182.8(4.9) \times 10^{-10} \tag{56}$$

so we may have new physics contribution explaining the 3.3 σ discrepancy,

$$\delta a_\mu \equiv a_\mu(\text{Exp}) - a_\mu(\text{SM}) = (26.1 \pm 8.0) \times 10^{-10}. \tag{57}$$

In the context of SUSY [80, 81], muon $g-2$ has the same Feynman diagram structure as the cLFV process $\mu \to e\gamma$. The crucial difference is that muon $g-2$ is flavor-conserving process, while $\mu \to e\gamma$ violates lepton flavor, L_μ and L_e. Therefore, δa_μ and $\text{Br}(\mu \to e\gamma)$ have a strong correlation[82],

$$\text{Br}(\mu \to e\gamma) \simeq 3 \times 10^{-5} \Big(\frac{\delta a_\mu^{\text{SUSY}}}{10^{-9}}\Big)^2 \Big(\frac{(m_L^2)_{12}}{m_{\text{SUSY}}^2}\Big)^2. \tag{58}$$

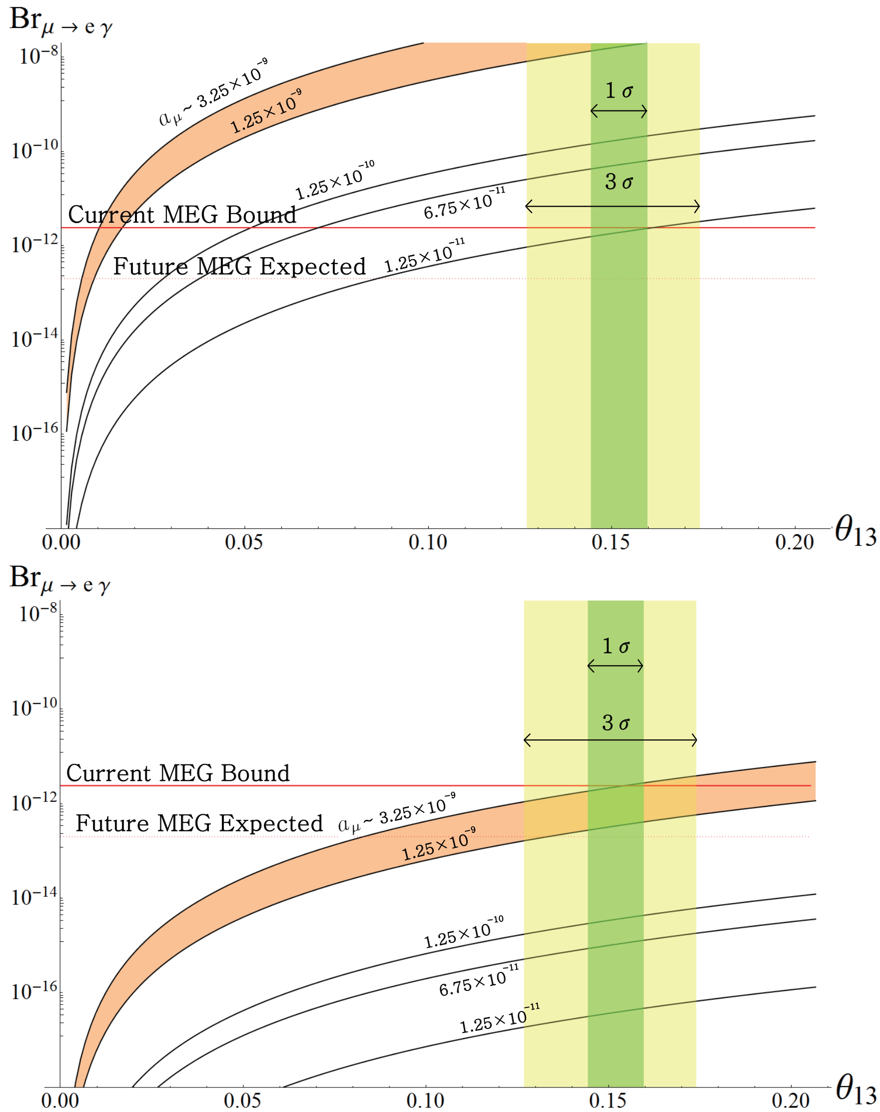

FIG. 14. Branching ratio of $\mu \to e\gamma$ as a function of θ_{13} for $\tan\beta = 10$, $y_\nu = 0.62$, $\rho = 0.1$. Future MEG expected bound is $O(10^{-13})$, we set the value 2×10^{-13}. Observed muon $g-2$ discrepancy is about $(2.25 \pm 1) \times 10^{-9}$, we draw $\mathrm{Br}(\mu \to e\gamma)$ at each muon $g-2$ contribution. Green and yellow band indicate 1σ, 3σ level of neutrino θ_{13}, respectively. In upper figure, θ_{13} is purely obtained from neutrino Dirac Yukawa splitting. In lower figure, only 1/15 portion of θ_{13} is obtained from neutrino Dirac Yukawa.

Moreover, the neutrino Dirac Yukawa Y_ν contains information on the neutrino oscillation observables. Since we consider the model where parameters of Y_ν are related to θ_{13} and cLFV, we have a strong correlation between $\mathrm{Br}(\mu \to e\gamma)$, θ_{13}, and muon $g-2$ as discussed

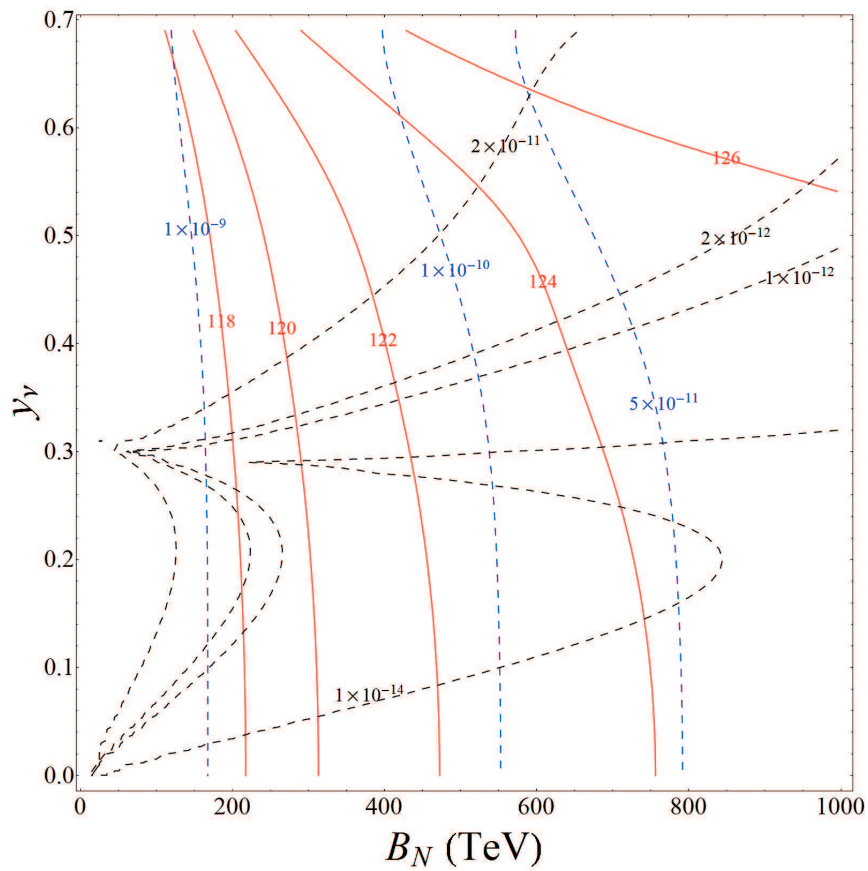

FIG. 15. Contour plot of Higgs mass(red solid line), cLFV(black dashed line), and muon $g-2$(blue dashed line) in B_N - y_ν plane for $\rho = 0.1$, $\tan\beta = 30$.

in [82].

Fig. 14 summarises the result. Both muon anomalous magnetic moment and cLFV is a function of $\tan\beta/M^2$ in which M is the typical supersymmetry breaking scale. The cLFV has extra suppression proportional to $(m_L^2)_{12}$. The S_4 flavor model discussed here is constructed from the neutrino Dirac Yukawa matrix which is proportional to the identity matrix and does not provide any off-diagonal entry in the slepton mass squared matrix if θ_{13} vanishes. Recently measured $\theta_{13} \sim 0.15$ provides an extra information depending on the origin of modification for nonzero θ_{13}.

If the full θ_{13} is explained by the degeneracy lift of the neutrino Dirac Yukawa matrix and if the entire discrepancy of the muon anomalous magnetic moment should be explained by light slepton, the current MEG bound tells that θ_{13} should be smaller than 0.01 which is incompatible with the observation recently made. The parameter space which is consistent

with $\mu \to e\gamma$ bound and the θ_{13} predicts that muon anomalous magnetic moment is at least 1/20 times smaller than what is needed.

If the nonzero θ_{13} is entirely generated by modifying the neutrino Majorana mass matrix, there would be no cLFV even for the sizeable θ_{13}. In reality, the subleading corrections in the simple flavor model would appear in both sectors and the observed θ_{13} would be a combined result of various sources. The bottom plot of Fig. 14 shows the hybrid case in which only 1/15 of the θ_{13} is from the neutrino Dirac Yukawa modification. In this case θ_{13}, muon anomalous magnetic moment can be explained at the same time. The $\mu \to e\gamma$ bound is satisfied and the the consistent region can be reached by the planned future MEG experiment as it predicts larger branching ratio of $\mu \to e\gamma$ than the planned expected sensitivity.

Fig. 15 shows the tension between the muon $g-2$ and the Higgs mass. Even if we take the model in which the neutrino Dirac Yukawa remains to be proportional to the identiy matrix such that no cLFV constraints apply, 125 GeV Higgs mass needs B_N much larger thna 300 TeV. Then the slepton is too heavy and the muon $g-2$ is much smaller than 10^{-9}. The figure also shows an interesting feature that the off-diagonal elements of m_L^2 vanish at around $y_\nu = 0.3$ and the cLFV bounds are very weak at around $y_\nu = 0.3$.

V. CONCLUSION

We considered the right-handed neutrinos as messengers of supersymmetry breaking in the minimal gauge mediation. Direct coupling of neutrino messenger with the Higgs field H_u and the lepton doublets L_i provides soft-trilinear A term for the top Yukawa and can help increase the light Higgs mass by realising the maximal stop mixing scenario. We call this setup as 'neutrino assisted gauge mediation'. The Yukawa mediation given by neutrino messengers also appear at soft scalar masses of the Higgs H_u, the lepton doublets L_i. At the same time it affects the soft scalar masses of the fields which couple to H_u and L_i at two loop. Among those, the stop mass squared gets the largest correction as the top Yukawa coupling is of order one. For y_ν slightly larger than 0.7, the correction is big enough to make stop tachyonic. Therefore, this realises the natural supersymmetry spectrum. At the same time maximal mixing is achieved by two effects, large A_t and small $m_{\tilde{t}}^2$ at around $y_\nu \sim 0.7$. In general this effect allows to explain the observed Higgs mass at around 125 GeV using around 1 TeV stop mass. Compared to the case when the neutrino assistance is turned off

$(y_\nu = 0)$, about 5 GeV of the Higgs mass is enhanced.

In general the off-diagonal entry of the slepton mass squared, m_L^2, appears at the messenger scale and can make the charged lepton flavor violating process to occur. The detailed quantitative prediction of cLFV highly depends on flavor model building. We provided a representative model based on S_4 flavor symmetry in which the Dirac neutrino Yukawa can be proportional to the identity if $\theta_{13} = 0$. For nonzero θ_{13}, two options were considered. Firstly, the total θ_{13} can be explained by the modification of the neutrino Dirac Yukawa matrix. Secondly, the θ_{13} can be explained by modifying the Majorana mass matrix of neutrinos. For the former, very stringent bound on the slepton mass comes from $\mu \to e\gamma$ and the slepton should be heavier than $2 \sim 4$ TeV, depending on $\tan\beta$. Also for the slepton mass at around 2 TeV with $\tan\beta = 10$, the $\mu \to e\gamma$ is just below the current experimental bound and we expect to observe the $\mu \to e\gamma$ in the near future.

Even for the second case in which we can safely avoid cLFV constraints, the neutrino assisted gauge mediation (in its minimal form with one copy of 5 and $\bar{5}$ messenger) sets the lower bound on the slepton mass to explain the Higgs mass. $1 \sim 2$ TeV slepton mass at the same time sets an upper bound on the possible contribution to muon anomalous magnetic moment and $a_\mu \sim 10^{-10}$ is the upper bound.

In this paper we proposed the neutrino assisted gauge mediation and showed a possible way out to avoid the strong cLFV constraints. Even then the current scheme has a tension with the muon anomalous magnetic moment which needs a lighter slepton. The extension of the minimal neutrino assisted gauge mediation to multiple messengers might ameliorate the tension between the spectrum needed to explain the Higgs mass and the muon anomalous magnetic moment.

ACKNOWLEDGMENTS

This work is supported by the NRF of Korea No. 2011-0017051.

	$(\tan\beta = 10, B_N = 360\,\mathrm{TeV})$	$(\tan\beta = 30, B_N = 300\,\mathrm{TeV})$
$\tilde{\nu}_e, \tilde{\nu}_\mu, \tilde{\nu}_\tau$	2957, 2961, 3013	2429, 2465, 2502
$\tilde{e}_1, \tilde{\mu}_1, \tilde{\tau}_1$	1364, 1364, 1333	1139, 1138, 880
$\tilde{e}_2, \tilde{\mu}_2, \tilde{\tau}_2$	3013, 2962, 2954	2503, 2467, 2427
$\tilde{u}_1, \tilde{c}_1, \tilde{t}_1$	2827, 2827, 634	2384, 2384, 637
$\tilde{d}_1, \tilde{s}_1, \tilde{b}_1$	2853, 2853, 2820	2406, 2406, 2283
$\tilde{u}_2, \tilde{c}_2, \tilde{t}_2$	3177, 3177, 2252	2675, 2675, 1868
$\tilde{d}_2, \tilde{s}_2, \tilde{b}_2$	3178, 3178, 2297	2676, 2676, 1893
$h_0, A, H_0, H_\pm$	125, 1705, 1705, 1707	125, 1031, 1031, 1034
$\chi_1, \chi_2, \chi_3, \chi_4$	487, 850, -892, 980	405, 713, -758, 829
χ_+, χ_-	849, 980	712, 829
$\tilde{g}$	2514	2126

TABLE III. Sparticle spectrum at the point giving 125 GeV Higgs mass with the lowest B_N

APPENDIX 0: SPARTICLE SPECTRUM SAMPLE POINT

APPENDIX A: REPRESENTATIONS OF S_4 SYMMETRY AND TENSOR PRODUCTS

S_4 is a non-abelian discrete symmetry and consists of all permutations among four quantities. For a review, see [83]. Irreducible representations of S_4 are two singlets $\mathbf{1}, \mathbf{1}'$, one singlet $\mathbf{2}$, and two triplets $\mathbf{3}, \mathbf{3}'$. Tensor products among them are given as follows:

$$\begin{aligned}\begin{pmatrix} x_1 \\ x_2 \\ x_3 \end{pmatrix}_{\mathbf{3}} \times \begin{pmatrix} y_1 \\ y_2 \\ y_3 \end{pmatrix}_{\mathbf{3}} &= (x_1 y_1 + x_2 y_2 + x_3 y_3)_{\mathbf{1}} + \begin{pmatrix} x_1 y_1 + \omega x_2 y_2 + \omega^2 x_3 y_3 \\ x_1 y_1 + \omega^2 x_2 y_2 + \omega x_3 y_3 \end{pmatrix}_{\mathbf{2}} \\ &+ \begin{pmatrix} x_2 y_3 + x_3 y_2 \\ x_3 y_1 + x_1 y_3 \\ x_1 y_2 + x_2 y_1 \end{pmatrix}_{\mathbf{3}} + \begin{pmatrix} x_2 y_3 - x_3 y_2 \\ x_3 y_1 - x_1 y_3 \\ x_1 y_2 - x_2 y_1 \end{pmatrix}_{\mathbf{3}'}\end{aligned} \tag{59}$$

	$(\tan\beta = 10, B_N = 240\,\mathrm{TeV})$	$(\tan\beta = 30, B_N = 200\,\mathrm{TeV})$
$\tilde{\nu}_e, \tilde{\nu}_\mu, \tilde{\nu}_\tau$	1971, 1974, 2009	1657, 1682, 1707
$\tilde{e}_1, \tilde{\mu}_1, \tilde{\tau}_1$	915, 915, 894	770, 769, 590
$\tilde{e}_2, \tilde{\mu}_2, \tilde{\tau}_2$	2010, 1976, 1970	1709, 1684, 1657
$\tilde{u}_1, \tilde{c}_1, \tilde{t}_1$	1937, 1937, 521	1633, 1633, 404
$\tilde{d}_1, \tilde{s}_1, \tilde{b}_1$	1954, 1954, 1931	1650, 1650, 1564
$\tilde{u}_2, \tilde{c}_2, \tilde{t}_2$	2169, 2169, 1569	1828, 1828, 1286
$\tilde{d}_2, \tilde{s}_2, \tilde{b}_2$	2170, 2170, 1586	1829, 1829, 1291
$h_0, A, H_0, H_\pm$	123, 1220, 1220, 1223	123, 679, 679, 684
$\chi_1, \chi_2, \chi_3, \chi_4$	322, 600, -729, 757	267, 466, -520, 579
χ_+, χ_-	600, 757	465, 579
$\tilde{g}$	1737	1470

TABLE IV. Sparticle spectrum at the point giving 123 GeV Higgs mass with the lowest B_N

$$\begin{pmatrix} x_1 \\ x_2 \\ x_3 \end{pmatrix}_{\mathbf{3'}} \times \begin{pmatrix} y_1 \\ y_2 \\ y_3 \end{pmatrix}_{\mathbf{3'}} = (x_1y_1 + x_2y_2 + x_3y_3)_{\mathbf{1}} + \begin{pmatrix} x_1y_1 + \omega x_2y_2 + \omega^2 x_3y_3 \\ x_1y_1 + \omega^2 x_2y_2 + \omega x_3y_3 \end{pmatrix}_{\mathbf{2}} + \begin{pmatrix} x_2y_3 + x_3y_2 \\ x_3y_1 + x_1y_3 \\ x_1y_2 + x_2y_1 \end{pmatrix}_{\mathbf{3}} + \begin{pmatrix} x_2y_3 - x_3y_2 \\ x_3y_1 - x_1y_3 \\ x_1y_2 - x_2y_1 \end{pmatrix}_{\mathbf{3'}} \tag{60}$$

$$\begin{pmatrix} x_1 \\ x_2 \\ x_3 \end{pmatrix}_{\mathbf{3}} \times \begin{pmatrix} y_1 \\ y_2 \\ y_3 \end{pmatrix}_{\mathbf{3'}} = (x_1y_1 + x_2y_2 + x_3y_3)_{\mathbf{1'}} + \begin{pmatrix} x_1y_1 + \omega x_2y_2 + \omega^2 x_3y_3 \\ -(x_1y_1 + \omega^2 x_2y_2 + \omega x_3y_3) \end{pmatrix}_{\mathbf{2}} + \begin{pmatrix} x_2y_3 + x_3y_2 \\ x_3y_1 + x_1y_3 \\ x_1y_2 + x_2y_1 \end{pmatrix}_{\mathbf{3'}} + \begin{pmatrix} x_2y_3 - x_3y_2 \\ x_3y_1 - x_1y_3 \\ x_1y_2 - x_2y_1 \end{pmatrix}_{\mathbf{3}} \tag{61}$$

$$\begin{pmatrix} x_1 \\ x_2 \end{pmatrix}_{\mathbf{2}} \times \begin{pmatrix} y_1 \\ y_2 \end{pmatrix}_{\mathbf{2}} = (x_1 y_2 + x_2 y_1)_{\mathbf{1}} + (x_1 y_2 - x_2 y_1)_{\mathbf{1}'} + \begin{pmatrix} x_2 y_2 \\ x_1 y_1 \end{pmatrix}_{\mathbf{2}} \tag{62}$$

$$\begin{pmatrix} x_1 \\ x_2 \end{pmatrix}_{\mathbf{2}} \times \begin{pmatrix} y_1 \\ y_2 \\ y_3 \end{pmatrix}_{\mathbf{3}} = \begin{pmatrix} (x_1 + x_2) y_1 \\ (\omega^2 x_1 + \omega x_2) y_2 \\ (\omega x_1 + \omega^2 x_2) y_3 \end{pmatrix}_{\mathbf{3}} + \begin{pmatrix} (x_1 - x_2) y_1 \\ (\omega^2 x_1 - \omega x_2) y_2 \\ (\omega x_1 - \omega^2 x_2) y_3 \end{pmatrix}_{\mathbf{3}'} \tag{63}$$

$$\begin{pmatrix} x_1 \\ x_2 \end{pmatrix}_{\mathbf{2}} \times \begin{pmatrix} y_1 \\ y_2 \\ y_3 \end{pmatrix}_{\mathbf{3}'} = \begin{pmatrix} (x_1 + x_2) y_1 \\ (\omega^2 x_1 + \omega x_2) y_2 \\ (\omega x_1 + \omega^2 x_2) y_3 \end{pmatrix}_{\mathbf{3}'} + \begin{pmatrix} (x_1 - x_2) y_1 \\ (\omega^2 x_1 - \omega x_2) y_2 \\ (\omega x_1 - \omega^2 x_2) y_3 \end{pmatrix}_{\mathbf{3}} \tag{64}$$

and trivially, we have $\mathbf{3} \times \mathbf{1}' = \mathbf{3}'$, $\mathbf{3}' \times \mathbf{1}' = \mathbf{3}'$, and $\mathbf{2} \times \mathbf{1}' = \mathbf{2}$.

APPENDIX B: REMARKS ON THE FLAVON VACUUM STABILITY

In [55], it was shown that A_4 triplet flavon vacuum in the direction of $(1,1,1)$ and $(1,0,0)$, ($(0,1,0)$, $(0,0,1)$ are the same) is favored compared to other directions, such as $(1,1,0)$. Since A_4 symmetry is the subgroup of the S_4 composed of even permutations, similar arguments hold. In this Appendix, we argue that triplet flavon directions favored in A_4 model are also favored in the S_4 model and that S_4 doublet vacuum favors the $(1,1)$ direction.

Rigid SUSY makes the discussion more simple, because the potential V is minimized at $\langle V \rangle = 0$. On the other hand, extra symmetries like Z_4 and $\mathrm{U}(1)_L$ more restrict possible terms in the superpotential. Suppose that $\mathrm{U}(1)_L$ symmetry is discretized to, for example, Z_8 symmetry. In this case, only quartic terms Φ^4 and χ^4 are allowed. Let us assume that breaking of extra symmetries introduces quadratic term, like $m_1 \Phi^2$ or $m_2 \chi^2$. To achieve this, let us consider 'Z_4 breaking singlets' ψ_1, $\bar{\psi}_1$ and 'lepton number breaking singlets' ψ_2, $\bar{\psi}_2$ with $S_4 \times Z_4 \times \mathrm{U}(1)_L$ quantum numbers

$$\begin{aligned} &\psi_1 : (\mathbf{1}, 3, 0), \quad \bar{\psi}_1 : (\mathbf{1}, 1, 0), \\ &\psi_2 : (\mathbf{1}, 0, 2), \quad \bar{\psi}_2 : (\mathbf{1}, 0, 6). \end{aligned} \tag{65}$$

They do not combine with $\bar{E}LH_d$, NLH_u and NN to make singlets under all symmetries imposed. Then, they can couple to Φ^2 or χ^2 such that a superpotential is given by

$$\begin{aligned}W(\psi_1,\bar{\psi}_1,\psi_2,\bar{\psi}_2) =& \frac{1}{\Lambda}[\Phi^2\bar{\psi}_1\psi_1+\chi^2\bar{\psi}_2\psi_2] \\ &- M_1\bar{\psi}_1\psi_1+\frac{1}{\Lambda}[\kappa_1(\bar{\psi}_1\psi_1)^2+\kappa_2(\psi_1)^4+\kappa_3(\bar{\psi}_1)^4] \\ &- M_2\bar{\psi}_2\psi_2+\frac{1}{\Lambda}[\kappa_1'(\bar{\psi}_2\psi_2)^2+\kappa_2'(\psi_2)^4+\kappa_3'(\bar{\psi}_2)^4].\end{aligned} \tag{66}$$

In this superpotential, $\bar{\psi}_1\psi_1$ and $\bar{\psi}_2\psi_2$ pairs have VEVs and they provide $m_1\Phi^2+m_2\chi^2$ terms. With this setup, the triplet superpotential has the form of

$$\begin{aligned}W = mS^2 &+ \frac{\lambda_1}{\Lambda}(x^2+y^2+z^2)^2+\frac{\lambda_2}{\Lambda}(x^2+\omega y^2+\omega^2z^2)(x^2+\omega^2y^2+\omega z^2) \\ &+ \frac{\lambda_3}{\Lambda}(xy+yz+zx)^2\end{aligned} \tag{67}$$

where $S=(x,y,z)$ represents the generic S_4 triplet such as Φ or χ. Note also that the superpotential has an accidental Z_2 symmetry under which $\psi_{1,2}$ and $\bar{\psi}_{1,2}$ are odd whereas other fields are even. If this Z_2 symmetry is imposed, $(\Phi^2\psi_1/\Lambda^3)\bar{E}LH_d$ and $(\Phi^2\psi_2/\Lambda^2)NN$ terms, which change the flavor structure in the subleading orders are forbidden. In this case, charged lepton Yukawa coupling structure in dimension-4 operator is preserved up to dimension-6 operator whereas Majorana mass structure in dimension-3 operator is preserved up to dimension-5 operator so corrections to them are highly suppressed.

Each term of the F-term potential $V=|\partial W/\partial x|^2+|\partial W/\partial y|^2+|\partial W/\partial z|^2$ is given by

$$\begin{aligned}\frac{\partial W}{\partial x} &= mx+\frac{4\lambda_1}{\Lambda}x(x^2+y^2+z^2)+\frac{2\lambda_2}{\Lambda}x(2x^2-y^2-z^2)+\frac{2\lambda_3}{\Lambda}(y+z)(xy+yz+zx) \\ \frac{\partial W}{\partial y} &= my+\frac{4\lambda_1}{\Lambda}x(x^2+y^2+z^2)+\frac{2\lambda_2}{\Lambda}y(2y^2-z^2-x^2)+\frac{2\lambda_3}{\Lambda}(z+x)(xy+yz+zx) \\ \frac{\partial W}{\partial z} &= mz+\frac{4\lambda_1}{\Lambda}z(x^2+y^2+z^2)+\frac{2\lambda_2}{\Lambda}z(2z^2-x^2-y^2)+\frac{2\lambda_3}{\Lambda}(x+y)(xy+yz+zx).\end{aligned} \tag{68}$$

Stable vacuum requires that these three terms should be zero simultaneously. For vacuum $\langle S\rangle = v(1,1,1)$, three terms give the same condition,

$$12(\lambda_1+\lambda_3)\Big(\frac{v^3}{\Lambda}\Big)+mv=0 \tag{69}$$

so the vacuum is stabilized at $v^2=-m\Lambda/[12(\lambda_1+\lambda_3)]$. For vacuum $\langle S\rangle = v(1,0,0)$, the second and third terms vanish trivially and the first term gives

$$4(\lambda_1+\lambda_3)\Big(\frac{v^3}{\Lambda}\Big)+mv=0 \tag{70}$$

so the vacuum is stabilized at $v^2 = -m\Lambda/[4(\lambda_1+\lambda_3)]$. The vacuum in the direction $(0,1,0)$ and $(0,0,1)$ gives the same result by permutational property of S_4. On the other hand, vacuum $\langle S\rangle = v(1,1,0)$ gives two conditions,

$$\begin{aligned}
&\frac{v^3}{\Lambda}(8\lambda_1+2\lambda_2+2\lambda_3)+mv=0\\
&\lambda_3 v^3 = 0.
\end{aligned} \tag{71}$$

If λ_3 is not forbidden by another symmetry, $v=0$ is the only solution and nontrivial vacuum can not be developed.

S_4 doublet stabilization can be discused in the same way. Renormalizable superpotential for doublet (x,y) is written as

$$W = m(xy)+\lambda(x^3+y^3) \tag{72}$$

and stabilization condition

$$\begin{aligned}
\frac{\partial W}{\partial x} &= 2my+3\lambda x^2 = 0\\
\frac{\partial W}{\partial y} &= 2mx+3\lambda y^2 = 0
\end{aligned} \tag{73}$$

requires that $x=y$. So the vacuum choice for Eq. (33) is stable.

APPENDIX C: COMMENT ON KÄHLER POTENTIAL CORRECTIONS

In our setup, Yukawa couplings are constructed from non-renormalizable dimension-4 superpotential with several flavons. These flavons also appear in the non-renormalizable Kähler potential and kinetic terms are written in the form of

$$K_{i\bar{j}}\partial_\mu\phi^{\bar{j}\dagger}\partial^\mu\phi^i - iK_{i\bar{j}}\bar{\psi}^{\bar{j}}\bar{\sigma}_\mu\partial^\mu\psi^i \tag{74}$$

where ϕ and ψ represent bosonic and fermionic fields, respectively. The Kähler potential of charged lepton supermultiplet L is given by

$$K = \left[1+a_1\frac{\Phi^\dagger\Phi}{\Lambda^2}+a_2\frac{\chi^\dagger\chi}{\Lambda^2}\right]L^\dagger L\Big|_{S_4\ \text{singlets}}+\cdots \tag{75}$$

and similar terms can be written for other fields, $\bar{E}^\dagger\bar{E}$, $N^\dagger N$, $H^\dagger_{u,d}H_{u,d}$, etc. Then we have quite complicate terms. For example, from $(\Phi^\dagger_{\mathbf{3}}\Phi_{\mathbf{3}}/\Lambda^2)L^\dagger L$ where $\Phi_{\mathbf{3}}$ vacuum is given by

$v_2(1,1,1)$, we have

$$a_1 \frac{\Phi_{\mathbf{3}}^\dagger \Phi_{\mathbf{3}}}{\Lambda^2} L^\dagger L\Big|_{S_4\ \text{singlets}} = a_{1,1}\frac{v_2^2}{\Lambda^2}(L_1^\dagger L_1 + L_2^\dagger L_2 + L_3^\dagger L_3) + a_{1,2}\frac{v_2^2}{\Lambda^2}\Big[L_2^\dagger L_3 + L_3^\dagger L_2 + L_3^\dagger L_1 + L_1^\dagger L_3 + L_1^\dagger L_2 + L_2^\dagger L_1\Big]. \tag{76}$$

Since $\langle\Phi\rangle/\Lambda = v_2/\Lambda$ is responsible for charged lepton Yukawa couplings, we see $4\pi v_2/\Lambda \gtrsim Y_\tau = m_\tau/[(v/\sqrt{2})\cos\beta] \sim 0.1$ for $\tan\beta = 10$. On the other hand, $\chi_{\mathbf{3}}$ has another vacuum direction, $w_2(0,1,0)$. Then

$$a_2 \frac{\chi^\dagger \chi}{\Lambda^2} L^\dagger L\Big|_{S_4\ \text{singlets}} = a_{2,1}\frac{w_2^2}{\Lambda^2}(L_1^\dagger L_1 + L_2^\dagger L_2 + L_3^\dagger L_3) + a_{2,2}\frac{w_2^2}{\Lambda^2}(-L_1^\dagger L_1 + L_2^\dagger L_2 - L_3^\dagger L_3) \tag{77}$$

so it just rescales the fields. Moreover, since See-Saw scale is about 10^{14} GeV, we have suppressed effect, $4\pi\chi/\Lambda \sim 0.01$ with Λ is the GUT scale. In the same way, doublet and singlet flavons in the Kähler potential just contribute to the field rescalings.

Physical fields are defined with canonical kinetic terms, so we should make field redefinitions and they affect flavor structures in principle. In our work, however, such effects are not considered by assuming small coeffecients $a_{1,2}$. For example, diagonalization of Y_E demonstrated above is not affected if $a_1(v_2^2/\Lambda^2) \lesssim (m_e/m_\tau) \sim 3\times 10^{-4}$, *i.e.* $a_1 \lesssim 3$.

On the other hand, mixings in the Kähler potential between flavons can be dangerous. For example, kinetic mixing between flavons such as $\bar{\psi}_1^\dagger \psi_2^\dagger \Phi_{\mathbf{3}}^\dagger \chi_{\mathbf{3}}/\Lambda^2$ can introduce small correction to Y_E or M_N with unwanted S_4 triplet vacuum direction. Such effect is suppressed by $\bar{\psi}_1^\dagger \psi_2^\dagger/\Lambda^2$ and can be more suppressed with tiny coefficient.

[1] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **716**, 1 (2012) [arXiv:1207.7214 [hep-ex]].

[2] S. Chatrchyan *et al.* [CMS Collaboration], Phys. Lett. B **716**, 30 (2012) [arXiv:1207.7235 [hep-ex]].

[3] M. Dine and A. E. Nelson, Phys. Rev. D **48**, 1277 (1993) [hep-ph/9303230].

[4] M. Dine, A. E. Nelson and Y. Shirman, Phys. Rev. D **51**, 1362 (1995) [hep-ph/9408384].

[5] M. Dine, A. E. Nelson, Y. Nir and Y. Shirman, Phys. Rev. D **53**, 2658 (1996) [hep-ph/9507378].

[6] G. F. Giudice and R. Rattazzi, Phys. Rept. **322**, 419 (1999) [hep-ph/9801271].

[7] M. A. Ajaib, I. Gogoladze, F. Nasir and Q. Shafi, Phys. Lett. B **713**, 462 (2012) [arXiv:1204.2856 [hep-ph]].

[8] J. L. Feng, Z. 'e. Surujon and H. -B. Yu, Phys. Rev. D **86**, 035003 (2012) [arXiv:1205.6480 [hep-ph]].

[9] K. J. Bae, K. Choi, E. J. Chun, S. H. Im, C. B. Park and C. S. Shin, arXiv:1208.2555 [hep-ph].

[10] S. P. Martin, Phys. Rev. D **81**, 035004 (2010) [arXiv:0910.2732 [hep-ph]].

[11] S. P. Martin and J. D. Wells, Phys. Rev. D **86**, 035017 (2012) [arXiv:1206.2956 [hep-ph]].

[12] K. J. Bae, T. H. Jung and H. D. Kim, arXiv:1208.3748 [hep-ph].

[13] Z. Kang, T. Li, T. Liu, C. Tong and J. M. Yang, arXiv:1203.2336 [hep-ph].

[14] N. Craig, S. Knapen, D. Shih and Y. Zhao, arXiv:1206.4086 [hep-ph].

[15] Y. Shadmi and P. Z. Szabo, JHEP **1206**, 124 (2012) [arXiv:1103.0292 [hep-ph]].

[16] A. Albaid, K. S. Babu and K. S. Babu, arXiv:1207.1014 [hep-ph].

[17] M. Abdullah, I. Galon, Y. Shadmi and Y. Shirman, arXiv:1209.4904 [hep-ph].

[18] J. L. Evans, M. Ibe, S. Shirai and T. T. Yanagida, Phys. Rev. D **85**, 095004 (2012) [arXiv:1201.2611 [hep-ph]].

[19] J. L. Evans, M. Ibe and T. T. Yanagida, Phys. Lett. B **705**, 342 (2011) [arXiv:1107.3006 [hep-ph]].

[20] M. Buican, P. Meade, N. Seiberg and D. Shih, JHEP **0903**, 016 (2009) [arXiv:0812.3668 [hep-ph]].

[21] R. Dermisek and H. D. Kim, Phys. Rev. Lett. **96**, 211803 (2006) [hep-ph/0601036].

[22] R. Dermisek, H. D. Kim and I. -W. Kim, JHEP **0610**, 001 (2006) [hep-ph/0607169].

[23] K. Choi, E. J. Chun, H. D. Kim, W. I. Park and C. S. Shin, Phys. Rev. D **83**, 123503 (2011) [arXiv:1102.2900 [hep-ph]].

[24] G. R. Dvali, G. F. Giudice and A. Pomarol, Nucl. Phys. B **478**, 31 (1996) [hep-ph/9603238].

[25] G. F. Giudice, H. D. Kim and R. Rattazzi, Phys. Lett. B **660**, 545 (2008) [arXiv:0711.4448 [hep-ph]].

[26] F. R. Joaquim and A. Rossi, Phys. Rev. Lett. **97**, 181801 (2006) [hep-ph/0604083].

[27] R. N. Mohapatra, N. Okada and H. -B. Yu, Phys. Rev. D **78**, 075011 (2008) [arXiv:0807.4524 [hep-ph]].

[28] P. Fileviez Perez, H. Iminniyaz, G. Rodrigo and S. Spinner, Phys. Rev. D **81**, 095013 (2010) [arXiv:0911.1360 [hep-ph]].

[29] F. Borzumati and A. Masiero, Phys. Rev. Lett. **57**, 961 (1986).

[30] M. Ciuchini, A. Masiero, P. Paradisi, L. Silvestrini, S. K. Vempati and O. Vives, Nucl. Phys. B **783**, 112 (2007) [hep-ph/0702144 [HEP-PH]].

[31] P. F. Harrison, D. H. Perkins and W. G. Scott, Phys. Lett. B **530**, 167 (2002) [hep-ph/0202074].

[32] Y. Lin, Nucl. Phys. B **824**, 95 (2010) [arXiv:0905.3534 [hep-ph]].

[33] H. Ishimori and E. Ma, Phys. Rev. D **86**, 045030 (2012) [arXiv:1205.0075 [hep-ph]].

[34] G. Altarelli, F. Feruglio, L. Merlo and E. Stamou, JHEP **1208**, 021 (2012) [arXiv:1205.4670 [hep-ph]].

[35] S. F. King, Phys. Lett. B **718**, 136 (2012) [arXiv:1205.0506 [hep-ph]].

[36] P. Minkowski, Phys. Lett. B **67**, 421 (1977).

[37] T. Yanagida, Conf. Proc. C **7902131**, 95 (1979).

[38] T. Yanagida, Prog. Theor. Phys. **64**, 1103 (1980).

[39] M. Gell-Mann, P. Ramond and R. Slansky, Conf. Proc. C **790927**, 315 (1979).

[40] R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44**, 912 (1980).

[41] G. F. Giudice, P. Paradisi and A. Strumia, Phys. Lett. B **694**, 26 (2010) [arXiv:1003.2388 [hep-ph]].

[42] G. F. Giudice and R. Rattazzi, Nucl. Phys. B **511**, 25 (1998) [hep-ph/9706540].

[43] N. Arkani-Hamed, G. F. Giudice, M. A. Luty and R. Rattazzi, Phys. Rev. D **58**, 115005 (1998) [hep-ph/9803290].

[44] Z. Chacko and E. Ponton, Phys. Rev. D **66**, 095004 (2002) [hep-ph/0112190].

[45] D. Grossman and Y. Nir, Phys. Rev. D **85**, 055004 (2012) [arXiv:1111.5751 [hep-ph]].

[46] G. F. Giudice and R. Rattazzi, Nucl. Phys. B **757**, 19 (2006) [hep-ph/0606105].

[47] Y. Abe *et al.* [DOUBLE-CHOOZ Collaboration], Phys. Rev. Lett. **108**, 131801 (2012) [arXiv:1112.6353 [hep-ex]].

[48] M. Hartz [T2K Collaboration], arXiv:1201.1846 [hep-ex].

[49] P. Adamson *et al.* [MINOS Collaboration], Phys. Rev. Lett. **108**, 191801 (2012) [arXiv:1202.2772 [hep-ex]].

[50] F. P. An *et al.* [DAYA-BAY Collaboration], Phys. Rev. Lett. **108**, 171803 (2012) [arXiv:1203.1669 [hep-ex]].

[51] J. K. Ahn *et al.* [RENO Collaboration], Phys. Rev. Lett. **108**, 191802 (2012) [arXiv:1204.0626

[hep-ex]].

[52] X. -G. He, Y. -Y. Keum and R. R. Volkas, JHEP **0604**, 039 (2006) [hep-ph/0601001].

[53] X. -G. He and A. Zee, Phys. Lett. B **645**, 427 (2007) [hep-ph/0607163].

[54] X. -G. He and A. Zee, Phys. Rev. D **84**, 053004 (2011) [arXiv:1106.4359 [hep-ph]].

[55] Y. BenTov, X. -G. He and A. Zee, arXiv:1208.1062 [hep-ph].

[56] F. Bazzocchi and L. Merlo, arXiv:1205.5135 [hep-ph].

[57] J. Beringer *et al.* [Particle Data Group Collaboration], Phys. Rev. D **86**, 010001 (2012).

[58] G. L. Fogli, E. Lisi, A. Marrone, D. Montanino, A. Palazzo and A. M. Rotunno, Phys. Rev. D **86**, 013012 (2012) [arXiv:1205.5254 [hep-ph]].

[59] M. C. Gonzalez-Garcia, M. Maltoni, J. Salvado and T. Schwetz, arXiv:1209.3023 [hep-ph].

[60] J. Hisano, T. Moroi, K. Tobe and M. Yamaguchi, Phys. Rev. D **53**, 2442 (1996) [hep-ph/9510309].

[61] E. Arganda and M. J. Herrero, Phys. Rev. D **73**, 055003 (2006) [hep-ph/0510405].

[62] J. L. Hewett, H. Weerts, R. Brock, J. N. Butler, B. C. K. Casey, J. Collar, A. de Govea and R. Essig *et al.*, arXiv:1205.2671 [hep-ex].

[63] J. Adam *et al.* [MEG Collaboration], Phys. Rev. Lett. **107**, 171801 (2011) [arXiv:1107.5547 [hep-ex]].

[64] B. Aubert *et al.* [BABAR Collaboration], Phys. Rev. Lett. **104**, 021802 (2010) [arXiv:0908.2381 [hep-ex]].

[65] U. Bellgardt *et al.* [SINDRUM Collaboration], Nucl. Phys. B **299**, 1 (1988).

[66] K. Hayasaka, K. Inami, Y. Miyazaki, K. Arinstein, V. Aulchenko, T. Aushev, A. M. Bakich and A. Bay *et al.*, Phys. Lett. B **687**, 139 (2010) [arXiv:1001.3221 [hep-ex]].

[67] C. Dohmen *et al.* [SINDRUM II. Collaboration], Phys. Lett. B **317**, 631 (1993).

[68] W. H. Bertl *et al.* [SINDRUM II Collaboration], Eur. Phys. J. C **47**, 337 (2006).

[69] B. O'Leary *et al.* [SuperB Collaboration], arXiv:1008.1541 [hep-ex].

[70] A. Blondel *et al.*, http://www.psi.ch/mu3e/DocumentsEN/LOI_Mu3e_PSI.pdf

[71] The PRIME working group, unpublished; LOI to J-PARC 50-GeV PS, LOI-25, http://www-ps.kek.jp/jhf-np/LOIlist/pdf/L25.pdf

[72] A. Abada, D. Das, A. Vicente and C. Weiland, JHEP **1209**, 015 (2012) [arXiv:1206.6497 [hep-ph]].

[73] J. L. Lopez, D. V. Nanopoulos and X. Wang, Phys. Rev. D **49**, 366 (1994) [hep-ph/9308336].

[74] U. Chattopadhyay and P. Nath, Phys. Rev. D **53**, 1648 (1996) [hep-ph/9507386].

[75] M. Hirsch, F. Staub and A. Vicente, Phys. Rev. D **85**, 113013 (2012) [arXiv:1202.1825 [hep-ph]].

[76] R. Fok and G. D. Kribs, Phys. Rev. D **82**, 035010 (2010) [arXiv:1004.0556 [hep-ph]].

[77] R. Kitano, M. Koike and Y. Okada, Phys. Rev. D **66**, 096002 (2002) [Erratum-ibid. D **76**, 059902 (2007)] [hep-ph/0203110].

[78] G. W. Bennett *et al.* [Muon G-2 Collaboration], Phys. Rev. D **73**, 072003 (2006) [hep-ex/0602035].

[79] K. Hagiwara, R. Liao, A. D. Martin, D. Nomura and T. Teubner, J. Phys. G **38**, 085003 (2011) [arXiv:1105.3149 [hep-ph]].

[80] T. Moroi, Phys. Rev. D **53**, 6565 (1996) [Erratum-ibid. D **56**, 4424 (1997)] [hep-ph/9512396].

[81] G. -C. Cho, K. Hagiwara, Y. Matsumoto and D. Nomura, JHEP **1111**, 068 (2011) [arXiv:1104.1769 [hep-ph]].

[82] J. Hisano and K. Tobe, Phys. Lett. B **510**, 197 (2001) [hep-ph/0102315].

[83] H. Ishimori, T. Kobayashi, H. Ohki, Y. Shimizu, H. Okada and M. Tanimoto, Prog. Theor. Phys. Suppl. **183**, 1 (2010) [arXiv:1003.3552 [hep-th]].